PHYSICAL REVIEW LETTERS

VOLUME 2 JUNE 15, 1959 NUMBER 12

THEORY OF THE ORIGIN OF THE VERY LOW-FREQUENCY RADIO EMISSIONS FROM THE EARTH'S EXOSPHERE*

J. W. MacArthur
Marlboro College, Marlboro, Vermont

(Received May 13, 1959)

In the theory propounded by Gallet[1] to explain dawn chorus, the electromagnetic waves are generated by interaction between the natural thermal noise present in the exosphere and the beams of high-speed protons incident upon the earth's magnetic field from the sun. This interaction, of the travelling-wave tube type, requires that the electromagnetic wave and the moving ions have essentially the same velocity over an extensive region of space in order for an appreciable amount of energy to be transferred from the beam to the wave. The following theory removes this objection, yet gives the same results.

A cloud of protons, impinging on the earth's magnetic field at high speed, will in general be deflected into a helical path about the lines of force, rotating with angular frequency $\omega_s = He/Mc$, where the proton mass $M = 1836\, m_e$. Thus a proton gyrofrequency is set up which is 1/1836 of the electron gyrofrequency

$$\omega_H = He/m_e c.$$

This rotating cloud of protons may be considered as a source of electromagnetic waves having angular source frequency ω_s and moving forward with velocity V. Hence a Doppler shift in frequency will result, with the radiated angular frequency given by

$$\omega = \omega_s/(1 - V/v_p), \tag{1}$$

where v_p, the phase velocity of the wave, is

$$v_p = c/n. \tag{2}$$

Now, the index of refraction, n, in the "whistler mode" of propagation through the exosphere, is known[2] to be given by

$$n^2 = 1 + \omega_p^2/(\omega\omega_H - \omega^2), \tag{3}$$

where $\omega_p = (4\pi Ne^2/m_e)^{1/2}$, the plasma angular frequency.

Combining (1), (2), and (3), there results

$$\omega^3(1 - V^2/c^2) - \omega^2[\omega_H - (V^2/c^2)\,\omega_H + 2\omega_s] + \omega[\omega_s^2 + 2\omega_s\omega_H + (V^2/c^2)\omega_p^2] - (\omega_s^2\omega_H) = 0.$$

If $V^2 << c^2$ and $\omega_s << \omega_H$, which are both very good approximations, this reduces to

$$\omega^3 - \omega^2(\omega_H) + \omega[2\omega_s\omega_H + (V^2/c^2)\omega_p^2] - (\omega_s^2\omega_H) = 0. \tag{4}$$

This is the basic equation of the theory presented here, but clearly approximate solutions are needed to clarify its meaning.

First, if it be assumed that $\omega_s^2 << \omega^2$, a condition easily met except for very low frequencies and quite near the earth, Eq. (4) simplifies to

$$\omega^2 - \omega(\omega_H) + 2\omega_s\omega_H + (V^2/c^2)\omega_p^2 = 0.$$

The quantity $8\omega_s/\omega_H$ occurring in the solution has magnitude 0.0043 and so can be ignored compared to 1. Thus

$$\omega \doteq \tfrac{1}{2}\omega_H[1 \pm (1 - 4V^2\omega_p^2/c^2\omega_H^2)^{1/2}]. \tag{5}$$

This is precisely Gallet's equation.

Next, if $\omega << \omega_H$, which is valid except far out and for quite high frequencies, Eq. (4) becomes

$$\omega^2 - \omega[2\omega_s + (V^2\omega_p{}^2/c^2\omega_H)] + \omega_s{}^2 = 0.$$

Hence

$$\omega = \omega_s + \frac{V^2\omega_p{}^2}{2c^2\omega_H} \pm \left(\frac{V^4\omega_p{}^4}{4c^4\omega_H{}^2} + \frac{V^2\omega_p{}^2\omega_s}{c^2\omega_H}\right)^{1/2}. \quad (6)$$

In the special case where $\omega_s << V^2\omega_p{}^2/2c^2\omega_H$, we find

$$\omega \doteq V^2\omega_p{}^2/c^2\omega_H, \quad (7)$$

which is identical to Gallet's equivalent result.

Thus it has been shown that Gallet's equations can be derived as limiting cases of a more general equation resulting from a quite different mechanism.

Several types of phenomena, such as the rising tones often associated with the "tails" of whistlers, are no doubt more easily understood in terms of the travelling-wave tube mechanism. However, it is likewise possible that the Doppler and proton-gyro concept can be made to cover cases not explainable by the former theory, and further work is progressing along these lines.

Evidently, then, the two processes are complementary rather than mutually exclusive, and the full explanation of all types of dawn chorus will entail a combination of the two theories.

*This research was generously supported by John B. Orr of Great Barrington, Massachusetts.

[1]R. M. Gallet, Proc. Inst. Radio Engrs. 47, 211 (1959).

[2]See, for example, J. A. Ratcliffe, The Magneto-Ionic Theory and Its Applications to the Ionosphere (Cambridge University Press, Cambridge, 1959).

SEARCH FOR AN ELECTRIC DIPOLE MOMENT OF THE ELECTRON*

D. F. Nelson, A. A. Schupp, R. W. Pidd, and H. R. Crane
The Harrison M. Randall Laboratory of Physics, The University of Michigan, Ann Arbor, Michigan
(Received May 25, 1959)

It has been pointed out[1,2] that the detection of an electric dipole moment (EDM) in an elementary particle would constitute a violation of both time-reversal invariance and space-reflection invariance. Experimental upper limits on the size of an EDM for several elementary particles have been found.[3-7] A summary of these limits is given in Table I. Garwin and Lederman[8] have remarked that the method used in searching for the muon EDM could be extended to higher accuracy for the electron. The anomalous magnetic dipole moment (MDM) experiment of Schupp, Pidd, and Crane[9] is an experiment of this general type. The purpose of this note is to show that this experiment does yield a new upper limit on the size of an EDM of the electron and to describe an extension of the same experiment that is expected to have a considerably greater sensitivity for detecting an EDM of the electron.

Table I. Experimental upper limits of the electric dipole moments of elementary particles.

Particle	Electric dipole moment	Reference
Electron	$\lesssim 10^{-13}$ cm $\times e$	3, 4
Positron	$\lesssim 8 \times 10^{-13}$ cm $\times e$	4
Proton	$\lesssim 10^{-13}$ cm $\times e$	5
Neutron	$\lesssim 3 \times 10^{-20}$ cm $\times e$	6
Muon	$\lesssim 2 \times 10^{-15}$ cm $\times e$	7

In the experiment on the anomalous MDM an unpolarized beam of electrons of 100-kev energy is sent parallel to a magnetic field. A portion of the beam that has been polarized by Mott scattering from a thin gold foil at almost 90° is allowed to execute a helical path in the magnetic field (see Fig. 1). Application of a small trapping field allows the retention of the electrons in the nearly homogeneous magnetic field for times up to several hundred microseconds. The beam then strikes a second thin gold foil. Two Geiger counters, placed at a scattering angle of 80° and in orientations parallel and antiparallel to the magnetic field, are used to observe the Mott asymmetry as a function of the time the electrons are trapped in the field. The ratio of the counts in these two counters has a periodic variation with trapping time. The frequency of this variation, the so-called difference frequency, is the magnitude of the angular velocity of spin pre-

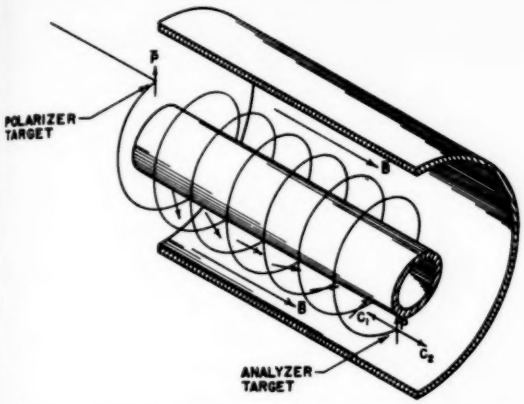

FIG. 1. Experimental geometry. The precession of the polarization vector $\vec{P}$ about the magnetic field $\vec{B}$ is indicated at several positions along the beam trajectory. The positions of the two counters are shown by C_1 and C_2.

cession relative to the velocity. In the absence of an EDM of the electron, this difference frequency is directly proportional to the anomalous part of the MDM. The existence of an EDM of the electron would increase this frequency by giving the electron spin an additional precession about the $\vec{v}\times\vec{B}$ direction. The contribution to this frequency of the EDM can, however, be distinguished by its energy dependence. Since in the anomalous MDM experiment the difference frequency was measured as a function of beam energy (the magnetic field being changed correspondingly to hold the beam at the same radius), a sensitive test of the existence of an EDM is contained in the result.

The general equation[10] governing the precession of the spin vector $\vec{\sigma}$ of an electron having an anomalous MDM and an EDM in both a magnetic field $\vec{B}$ and an electric field $\vec{E}$ (which vary only on a macroscopic scale[11]) is

$$d\vec{\sigma}/dt = \vec{\omega}_s \times \vec{\sigma}, \tag{1}$$

where

$$\vec{\omega}_s = -\frac{e}{m_0\gamma}\left[\vec{B} - \frac{\gamma}{\gamma+1}\frac{\vec{v}\times\vec{E}}{c^2}\right] - \frac{ea}{m_0}\left[\vec{B} - \frac{\gamma-1}{\gamma}\frac{\vec{B}\cdot\vec{v}}{v^2}\vec{v} - \frac{\vec{v}\times\vec{E}}{c^2}\right] - \frac{ef}{m_0 c}\left[\vec{E} - \frac{\gamma-1}{\gamma}\frac{\vec{E}\cdot\vec{v}}{v^2}\vec{v} + \vec{v}\times\vec{B}\right]. \tag{2}$$

The factors a and f are defined in terms of the MDM vector $\vec{\mu}$ and the EDM vector $\vec{p}$ by $\vec{\mu} = (1+a) \times e\hbar\vec{\sigma}/2m_0$, and $\vec{p} = fe\hbar\vec{\sigma}/2m_0c$; e, m_0, and $\vec{v}$ are the charge, rest mass, and velocity of the electron; $\gamma = (1-\beta^2)^{-1/2}$ and $\beta = v/c$. The difference frequency is defined as

$$\omega_D = |\vec{\omega}_s - \vec{\omega}_c|, \tag{3}$$

where $\vec{\omega}_c$ is the angular velocity of cyclotron rotation. For the case of $\vec{E}=0$ and $\vec{v}\cdot\vec{B}=0$

$$\omega_D = \omega_0(a^2+f^2\beta^2)^{\frac{1}{2}}, \tag{4}$$

where $\omega_0 = eB/m_0$. (The error introduced by assuming $\vec{v}\cdot\vec{B}=0$ is negligible.) It is seen that ω_D increases with increasing energy and that at infinite energy $\omega_D = (a^2+f^2)^{1/2}$. In the experiment on the anomalous MDM, $\omega_D{}^{\exp} > a$ when extrapolated to infinite energy, but $\omega_D{}^{\exp}$ decreases with increasing energy. (See Fig. 2.) The most likely

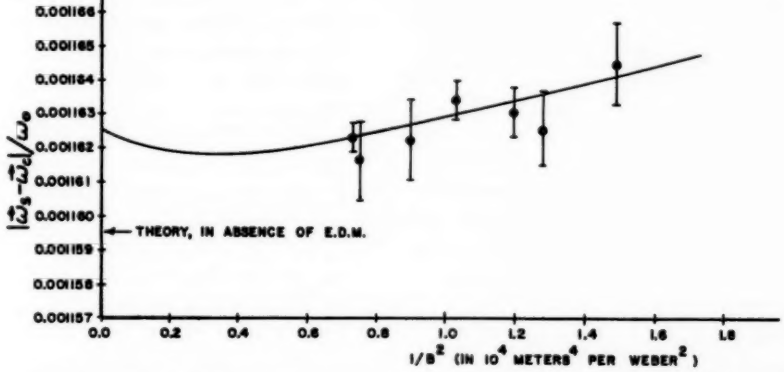

FIG. 2. Plot of difference frequency data. The experimental values of the difference frequency divided by eB/m_0 are extrapolated to infinite magnetic field (infinite energy). From the intercept $(a^2+f^2)^{1/2}$ the factor f (the EDM in units of $e\hbar/2m_0c$) can be found assuming the theoretical value of the g-factor anomaly a.

cause of this reversed slope is the presence of a small stray radial electric field. An electric field in any other direction, large enough to explain the slope, would have to be about 100 times larger than one that could reasonably be expected to exist in the apparatus. If one represents the stray electric field as a small radial electric field E, constant in magnitude at the orbit radius, ω_D is given by

$$\omega_D = \omega_0 \left\{ \left[a + \frac{E}{cB}\left(\frac{\beta^2 + a\beta^2 - 1}{\beta}\right)\right]^2 + f^2\left[\beta + \frac{E}{cB}\right]^2 \right\}^{\frac{1}{2}}. \quad (5)$$

If a two-parameter (f and E) fit of Eq. (5) is made to the data, the theoretical value of a being assumed, the results are

$$|p| = (1.6 \pm 0.6) \times 10^{-15} \text{ cm} \times e,$$

$$E = -(0.048 \pm 0.013) \text{ v/cm}. \quad (6)$$

It is reasonable that a stray electric field of this size existed in the apparatus. The errors quoted are standard deviations of statistical error. Systematic errors are estimated to be of the order of 0.15% in each measurement of ω_D/ω_0. When these errors are included, the results are consistent with the nonexistence of an EDM of the electron. It is concluded that

$$|p| \leq 3 \times 10^{-15} \text{ cm} \times e. \quad (7)$$

It appears that the smallest EDM detectable by this method is about 3×10^{-16} cm $\times e$. This would require more complete elimination of stray electric fields (and measurement of them, if possible), improvement in the accuracy of the measurement of ω_D/ω_0, and extension of the measurements of ω_D/ω_0 to higher energy (~600 kev).

The fundamental difficulty of this method of searching for an EDM is that the precession of the spin produced by the EDM is small compared with that produced by the anomalous part of the MDM. This difficulty can be removed if the magnetic field in the system of the electron can be reduced enough so that the spin precession frequency due to the whole MDM equals the cyclotron rotation frequency. This can be done by providing a small radial electric field (~20 v/cm). In the absence of an EDM the effect of the radial electric field is to fix the spin vector relative to the velocity vector. If there is an EDM there will be a precession of the spin caused by the EDM around the radial direction ($\vec{v} \times \vec{B}$), and there will not be any anomalous MDM precession present to obscure it.

For the radial electric field scheme to work, the polarization vector of the beam cannot initially be in the radial direction since there would then be no torque on the EDM. Since this is the initial orientation of the polarization vector in the experiment on the anomalous MDM, a modification of this setup is necessary. The modification is to allow the polarization vector to precess, due to the anomalous MDM, to an orientation parallel to the velocity before switching on the radial electric field. The only further modification needed is the re-orientation of the counters in the radial direction. The ratio of the counts in the two counters will then vary with trapping time according to the frequency of the precession produced by the EDM, given by

$$\omega_D = \omega_0 f(\beta + E/cB), \quad (8)$$

where the electric field is determined by

$$a + (E/cB)[(\beta^2 + a\beta^2 - 1)/\beta] = 0. \quad (9)$$

The amplitude of this variation will be the same as that observed in the anomalous MDM experiment. A second pair of counters, oriented like those in the anomalous MDM experiment, will be used to determine experimentally the proper adjustment of the radial electric field. The absence of a variation in the ratio of counts in these two counters when the trapping time is changed will indicate that the electric field is properly adjusted.

The practical limit of sensitivity of this method is imposed by the depolarization arising from the energy width of the beam in conjunction with the variation of the electric field with radius. Electrons of different energy will occupy orbits of different radii and so experience different electric fields. At only one radius will the electric field be such as to eliminate the anomalous MDM precession exactly. At other radii the anomalous MDM will cause additional precession which will cause a depolarization. If a $1/r$ falloff of the electric field and an energy width of 100 ev at 100-kev beam energy (caused by ionization losses in the first scattering foil) are assumed, a useful trapping time of 3×10^{-3} sec results. Since the phase of the EDM precession can be changed at will by changing the time the radial electric field is switched on, a tenth of a period of the EDM precession can be distinguished from spurious effects. An upper limit on the EDM of the electron of about 4×10^{-18} cm $\times e$ would result. When working to this accuracy, effects due to the component of the velocity parallel to the mag-

netic field, to stray electric fields, and to radial magnetic fields in the trapping region can be made negligibly small.

We would like to thank Dr. L. M. Lederman for a stimulating communication.

---

*This work is supported by the U. S. Atomic Energy Commission.

[1]L. Landau, Nuclear Phys. 3, 127 (1957).
[2]T. D. Lee and C. N. Yang, Brookhaven National Laboratory Report BNL-443, 1957 (unpublished).
[3]G. Feinberg, Phys. Rev. 112, 1637 (1958).
[4]E. E. Salpeter, Phys. Rev. 112, 1642 (1958).
[5]R. M. Sternheimer, Phys. Rev. 113, 828 (1959).
[6]Smith, Purcell, and Ramsey, Phys. Rev. 108, 120 (1957).
[7]Berley, Garwin, Gidal, and Lederman, Phys. Rev. Letters 1, 144 (1958).
[8]R. L. Garwin and L. M. Lederman, Nuovo cimento (to be published).
[9]Schupp, Pidd, and Crane, Bull. Am. Phys. Soc. Ser. II, 4, 250 (1959); and to be published.
[10]G. W. Ford (unpublished).
[11]These equations refer to expectation values of the various operators over a wave packet whose dimensions are small compared with the lengths in which the fields vary appreciably.

THEORY OF THE LATTICE VIBRATIONS OF GERMANIUM

W. Cochran*
Atomic Energy of Canada Limited, Chalk River, Ontario, Canada
(Received May 15, 1959)

The relation between frequency ω and wave vector $\vec{q}$ (the dispersion relation) for the normal modes of germanium has been investigated experimentally by Brockhouse and Iyengar[1] (see also Brockhouse[2]) and by Ghose, Palevsky, Hughes, Pelah, and Eisenhauer.[3] The results do not agree with prior calculations made by Hsieh,[4] on the assumption that only nearest neighbor atoms interact and using the Born-von Kármán theory of lattice dynamics. Nor is it sufficient to introduce second neighbor interaction[1]; indeed Herman[5] has shown that force constants between atoms out to at least fifth neighbors are required to account for the experimental results. The theory then involves numerous parameters which have no clear physical significance. Lax[6] proposes to fit the data to a force model involving one parameter to represent electrostatic interaction between quadrupoles generated by the lattice vibrations, and as many near neighbor parameters as proves necessary.

We have extended the Born-von Kármán theory to apply to a model of the germanium crystal in which each atom is treated as a charged core coupled isotropically to an oppositely charged shell of negligible mass. This atom model has had some success in accounting for the dielectric properties[7,8] and lattice dynamics[9] of alkali halide crystals. Its use introduces two long-range forces, firstly the electrostatic interaction between dipoles formed by the relative displacement of cores and shells, and secondly a force which originates from the massless character of the shells. The shell model therefore simulates two effects which must occur in practice, electrostatic interaction and relatively widespread redistribution of valence electron density when the nuclei are displaced. The condition that the crystal as a whole can have no dipole moment[6] is automatically satisfied, so that the electrostatic interaction need not be treated as quadrupole-quadrupole.

In the Born-von Kármán theory the dispersion relation is determined by the condition for solubility of the equations

$$\sum_{k'y}[M_{xy}(kk') - \omega^2 m_k \delta_{kk'}]U_y(k') = 0, \qquad (1)$$

where $\vec{U}(k)\exp i[\vec{q}\cdot\vec{r}(lk) - \omega t]$ is the displacement from $\vec{r}(lk)$ of the kth atom in the lth cell. We define $M_{xy}(kk')$ by writing

$$M_{xy}(kk') = B_{xy}(kk') + C_{xy}(kk'),$$

where $B_{xy}(kk')$ is related to the bonding force constant $\Phi_{xy}(lk, l'k')$ between atoms (lk) and $(l'k')$ by the equation

$$B_{xy}(kk') = -\sum_{l'}\Phi_{xy}(lk, l'k') \times \exp i\vec{q}\cdot[\vec{r}(l'k') - \vec{r}(lk)], \qquad (2)$$

and $C_{xy}(kk')$ is related to the electrostatic interaction between charged atoms by a similar expression.[10] For a diamond-type crystal with $\vec{q}$ along

[100] or [111] the set of Eqs. (1) can be reduced to three separate sets, each of which determines the dispersion relation for a single mode, and which can be written

$$\sum_{k'}[M(kk') - \omega^2 m_k \delta_{kk'}]U(k') = 0 \qquad (3)$$

where in general M is a linear combination of M_{xx}, M_{xy}, etc. For the shell model $k = 1 \ldots 4$, indices 1 and 2 referring to the two cores per unit cell, indices 3 and 4 to the corresponding shells, and we take $m_1 = m_2 = m$, $m_3 = m_4 = 0$. The number of independent $\Phi_{xy}(kk')$ when only nearest neighbor units interact is six. It is convenient to define new coefficients

$$R = B(12) + B(32) + B(14) + B(34),$$

$$T = B(32) + B(34),$$

$$S = B(32).$$

Each is a function of $\vec{q}$ and of two parameters. We introduce the abbreviations $\rho = R/R_0$, $\tau = T/T_0$, and $\sigma = S/S_0$, where $(R)_{q=0} = -R_0$, and similarly for T_0 and S_0. We also abbreviate $\beta = \alpha/v$, where α is the atomic polarizability and v the volume of the (trigonal) unit cell. The dispersion relation derived from the set of Eqs. (3) is then

$$\begin{vmatrix} A_0 - m\omega^2 & A \\ A^* & A_0 - m\omega^2 \end{vmatrix} = 0, \qquad (4)$$

where

$$A_0 = R_0 + B_1\left\{\frac{-(1 + B_2 + \beta C_1)(1 + |\tau|^2) + \tau(\beta C_2^* + B_2\sigma^*) + \tau^*(\beta C_2 + B_2\sigma)}{(1 + B_2 + \beta C_1)^2 - |\beta C_2 + B_2\sigma|^2}\right\}, \qquad (5)$$

$$A = R + B_1\left\{\frac{-2\tau(1 + B_2 + \beta C_1) + (\beta C_2 + B_2\sigma) + \tau^2(\beta C_2^* + B_2\sigma^*)}{(1 + B_2 + \beta C_1)^2 - |\beta C_2 + B_2\sigma|^2}\right\}. \qquad (6)$$

Here B_1 and B_2 are constants, while C_1 and C_2 are dimensionless functions of $\vec{q}$ depending only on the geometry of the crystal structure. For example,

$$(C_1)_{xy} = -v \lim_{\vec{r}\to 0}\left[\sum_l{}' \frac{\partial^2}{\partial x \partial y}\left(\frac{1}{|\vec{r} - \vec{r}_l|}\right)\exp i\vec{q}\cdot\vec{r}_l\right],$$

where $\vec{r}_l$ is a lattice translation. These functions were evaluated by the Ewald method, as developed by Kellerman.[10] The fact that the elastic constants of germanium satisfy Born's[11] identity now appears as a reflection of the fact that there are bonding interactions only between nearest neighbor units (cores or shells), and that $\rho = \tau$. Conversely, the fact that the elastic constants of silicon and diamond do not satisfy the identity may merely indicate $\rho \neq \tau$, and not that there are other than first neighbor bonding interactions. Three of the seven parameters involved in the dispersion relation are determined by the elastic constants. On physical grounds we also set $\rho = \sigma$, thereby eliminating a fourth parameter and making $\Phi_{xy}(kk')/\Phi_{xx}(kk')$ the same for $kk' = 12$, 14, and 34. The value of β deduced from the dielectric constant of germanium[12] is 0.100, leaving only two free parameters, the constants B_1 and B_2. In fact, we also treated β as a disposable parameter in fitting the data of Brockhouse and Iyengar, and obtained $\beta = 0.105$, in satisfactory agreement with the value given by the dielectric constant. With the following numerical values,

$$R_0 = 2.923 \times 10^5 \text{ dyne cm}^{-1},$$

$$B_1 = 4.099 \times 10^5 \text{ dyne cm}^{-1},$$

$$B_2 = 3.200,$$

$$\beta = 0.105,$$

and

$$\Phi_{xy}/\Phi_{xx} = 0.6905 \quad \text{(in all three instances)}$$

we obtained a good fit to the experimental data, as is shown in Table I and Fig. 1. The chief de-

Table I. The elastic constants.

	C_{11}	C_{12}	C_{44}
Calculated	1.292	0.492	0.676
Measured,[a] 300°K	1.288	0.483	0.671
Measured,[b] 0°K	1.310	0.490	0.687

[a] H. J. McSkimin, J. Appl. Phys. 24, 988 (1953).
[b] M. E. Fine, J. Appl. Phys. 26, 862 (1955).

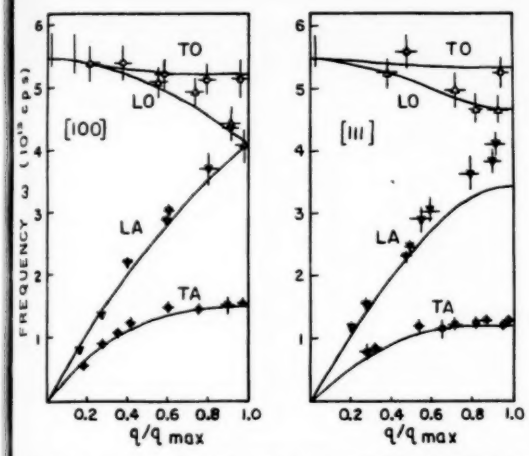

FIG. 1. A comparison of the theoretical dispersion curves with the experimental results of Brockhouse and Iyengar.

fect is that the calculated frequency of the LA mode is too low when $\vec{q}$ is near the zone boundary in the [111] direction.

It is remarkable that such a simple model should fit the data as well as it does. Our procedure, however, finds theoretical backing in the work of Mashkevich and Tolpygo.[13] Using the Hartree-Fock and adiabatic approximations, they have shown that the energy perturbation should be expressed as a quadratic function of the nuclear displacements and the atomic dipole moments. It can be shown that the shell model represents the most general (and probably the simplest) way of accomplishing this.

I am greatly indebted to Dr. B. N. Brockhouse for many helpful discussions and suggestions.

*On leave from Crystallographic Laboratory, Cavendish Laboratory, Cambridge, England.

[1]B. N. Brockhouse and P. K. Iyengar, Phys. Rev. 111, 747 (1958).

[2]B. N. Brockhouse, Phys. Rev. Letters 2, 256 (1959).

[3]Ghose, Palevsky, Hughes, Pelah, and Eisenhauer, Phys. Rev. 113, 49 (1959).

[4]Y. C. Hsieh, J. Chem. Phys. 22, 306 (1954).

[5]F. Herman, J. Phys. Chem. Solids, 8, 405 (1959). A similar calculation has been made by N. Pope (private communication).

[6]M. Lax, Phys. Rev. Letters 1, 133 (1958).

[7]B. G. Dick and A. W. Overhauser, Phys. Rev. 112, 90 (1958).

[8]J. E. Hanlon and A. W. Lawson, Phys. Rev. 113, 472 (1959).

[9]Woods, Cochran, and Brockhouse, Bull. Am. Phys. Soc. Ser. II, 4, 246 (1959).

[10]E. W. Kellerman, Trans. Roy. Soc. (London) 238, 513 (1940).

[11]M. Born, Ann. Physik 44, 605 (1914).

[12]H. Y. Fan, in Solid-State Physics, edited by F. Seitz and D. Turnbull (Academic Press, New York, 1955), Vol. 1, p. 281.

[13]V. S. Mashkevich and K. B. Tolpygo, J. Exptl. Theoret. Phys. U.S.S.R. 32, 520 (1957) [translation: Soviet Phys. JETP 5, 435 (1957)].

DETECTION OF COLLOIDAL CENTERS IN LITHIUM HYDRIDE BY ELECTRON RESONANCE

W. T. Doyle,* D. J. E. Ingram, and M. J. A. Smith
Department of Electronics, Southampton University, Southampton, England
(Received May 18, 1959)

This Letter reports electron resonance measurements on lithium hydride which has been irradiated with ultraviolet light and the results show that the lithium hydride contains colloidal centers of lithium metal, with diameters small compared with the microwave skin depth. The lithium hydride was irradiated for about fifteen minutes with 2537A radiation from a 250-watt mercury lamp and a fairly dark coloration was produced. These samples were then inserted in a 3-cm wavelength electron resonance spectrometer employing 100-kc/sec field modulation and phase-sensitive detection, and their electron resonance spectra recorded. All the irradiated samples gave a single resonance line at $g=2.002 \pm 0.001$ with a half-width of 0.3 gauss—the latter being determined by the magnetic field inhomogeneity. In contrast to this, no signals at all were obtained from the unirradiated material.

This very narrow signal shows that the paramagnetic entity cannot be due to normal electron excess centers, which have wide inhomogeneously broadened lines.[1] It was therefore postulated that the absorption was in fact due to the conduction electrons in colloidal particles of lithium. In order to test this hypothesis a careful inves-

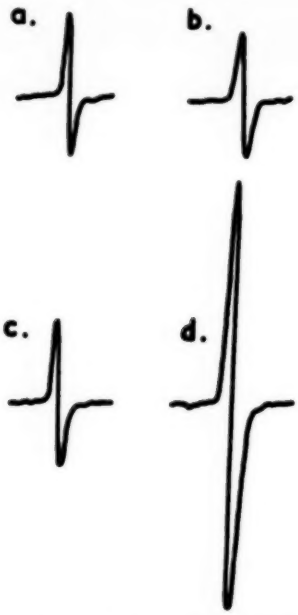

FIG. 1. Electron spin resonance of irradiated LiH. Curves (a) and (b) show the derivative of the absorption signals obtained from LiH and from a standard free-radical sample, respectively, at room temperature. Curves (c) and (d) show the corresponding signals at liquid oxygen temperature.

tigation on the temperature dependence of the intensity of the electron resonance was made, and the results are summarized in Fig. 1. Measurements were made at room temperature and at 90°K, and in each case the curves from the lithium sample and a standard free-radical sample were recorded under identical conditions. It can be seen from the figure that, whereas the intensity of the signal from the free radical increases markedly as the temperature is lowered, as would be expected, that from the irradiated lithium hydride remains constant. The fact that narrow symmetrical electron resonance lines are obtained with intensities independent of temperature shows that they must originate from colloidal particles of lithium with diameters small compared with the skin depth at microwave frequencies.[2]

Further confirmation of the presence of such colloidal particles is obtained from the optical absorption spectra. Optical absorption measurements have been made on a very small single crystal and the differential absorption of the irradiated and nonirradiated crystal is plotted in

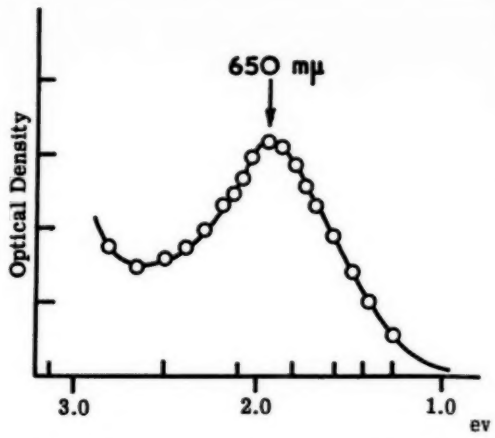

FIG. 2. Differential optical absorption spectrum of irradiated LiH. Optical density is plotted against photon energy in ev.

Fig. 2. It is evident that the irradiation has produced a strong band at 650 mμ. The position of this band does not agree at all with the positions predicted for F or aggregate bands in the hydrides,[3] whereas a strong band at 2512 A has been reported[4] and its position agrees well with that predicted for the F band. Preliminary estimates on the expected position of the colloid absorption band, based on the index of refraction of the host medium[5] and the plasma wavelength[6] of the alkali metal, can be made, however. These preliminary calculations suggest that, although very small colloidal particles would be expected to give rise to an absorption at a considerably shorter wavelength than 650 mμ, it is nevertheless quite possible for the band to be shifted to this region if larger colloidal particles are involved. The electron resonance measurements show that at least 10^{16} unpaired electrons per gram are present in the specimens, and hence a relatively large colloidal content of lithium is present.

The general conclusion is that both the electron resonance and optical absorption spectra show that colloidal centers of metallic lithium are formed when lithium hydride is irradiated with ultraviolet light. It would appear that this may be a very general effect, and initial studies on the thermal and optical decomposition of various azides have also yielded narrow electron spin resonance lines. Further work on these and other centers, employing much more homogeneous magnetic fields, is now in hand so that the

theories relating the width of the lines, the spin-orbit coupling, and the mean free path of the conduction electrons, can be checked in detail.

---

* National Science Foundation postdoctoral fellow. On leave from Dartmouth College, Hanover, New Hampshire.

[1]A. M. Portis, Phys. Rev. 91, 1071 (1953).
[2]F. J. Dyson, Phys. Rev. 98, 349 (1955).
[3]W. T. Doyle, Proc. Phys. Soc. (London) (to be published).
[4]F. Bach and K. F. Bonhoffer, Z. physik. Chem. 23B, 256 (1933); W. Rauch, Z. Physik 111, 650 (1938).
[5]H. Bode, Z. physik. Chem. 6B, 251 (1930).
[6]R. W. Wood and C. Lukens, Phys. Rev. 54, 332 (1938).

EFFECTIVE EXCHANGE CONSTANT IN YTTRIUM IRON GARNET*

D. T. Edmonds and R. G. Petersen
Department of Physics, University of California, Berkeley, California
(Received May 18, 1959)

We have measured the specific heat of a pure (99.99% yttrium) polycrystalline sample of yttrium iron garnet (YIG) in the temperature range of 1°K to 4°K. We find that the specific heat may be analyzed into a part that depends on temperature as $T^{3/2}$ (the spin-wave contribution) and a part that depends on temperature as T^3 (the contribution from the lattice) as illustrated in Fig. 1.

We thus obtain

$$C_v = 68.9T^{3/2} + 29.3T^3 \text{ erg/cm}^3 \text{ deg K}, \qquad (1)$$

where C_v is the specific heat per cm³ based on the x-ray density[1] of 5.19 g/cm³.

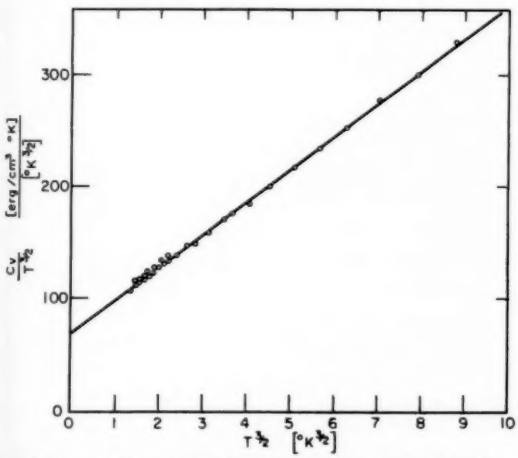

FIG. 1. The specific heat per unit volume of polycrystalline YIG analyzed into its two components by plotting $C_v/T^{3/2}$ against $T^{3/2}$. The points give the results obtained in one experiment. Another experiment performed on the same specimen at the extremes of the temperature range yielded results in agreement with those plotted to within 2%.

Using the equation

$$C_v(\text{lattice}) = (12/5)\pi^4 N k_B (T/\theta)^3, \qquad (2)$$

where k_B is the Boltzmann constant, we obtain for the Debye temperature $\theta = 454$°K, if we take N to be the total number of atoms of all kinds in unit volume.

Assuming a dispersion relation for the spin waves of

$$\hbar\omega = Dk^2, \quad k = 2\pi/\lambda, \qquad (3)$$

one has the standard result for the magnon specific heat per unit volume:

$$C_v(\text{magnons}) = [(15/32)(1.341)\pi^{-3/2}] k_B (k_B T/D)^{3/2}$$
$$= 0.113\, k_B (k_B T/D)^{3/2}. \qquad (4)$$

The constant D may be expressed[2] in terms of the Landau-Lifshitz exchange constant A, the saturation magnetization M_s, and the spectroscopic splitting factor g as

$$D = 2Ag(\mu_B/M_s), \qquad (5)$$

where μ_B is the Bohr magneton. In this manner we obtain from our measurement of the spin-wave contribution to the specific heat,

$$A(\text{YIG}) = 0.192 \times 10^{-6} \text{ erg/cm}, \qquad (6)$$

where we have taken the value[3] $4\pi M_s = 1740$ gauss for YIG. This is to be compared with a value of

$$A(\text{Fe}) \cong 2 \times 10^{-6} \text{ erg/cm}$$

in iron. If we analyze in this manner the results of Kouvel[4] who measured the specific heat of

magnetite in the same temperature range, we obtain

$$A(Fe_3O_4) = 0.232 \times 10^{-6} \text{ erg/cm}.$$

Suhl[5] has shown that in a sphere, z-directed spin waves having a wave number given by k_0 in

$$k_0^2 = 2\pi M_S^2/3A$$

are degenerate in frequency with the uniform precessional mode ($k = 0$). We expect[6,7] maximum coupling of the uniform precessional mode to transverse phonons via the z-directed spin waves when the frequency is given by

$$f_c = v_t k_0/2\pi,$$

where v_t is the velocity of transverse phonons. Using our value of A and the value $v_t = 3.87 \times 10^5$ cm/sec obtained in polycrystalline YIG at 25°C by McSkimin,[8] we find

$$|k_0| = 4.60 \times 10^5 \text{ cm}^{-1}$$

and

$$f_c = 2.83 \times 10^{10} \text{ sec}^{-1} \quad \text{or} \quad \lambda_c = 1.05 \text{ cm}.$$

It is our pleasure here to record our debt to Professor C. Kittel who suggested the investigation and to M. K. Jack of Hughes Research Laboratories who kindly supplied us with the specimen of YIG.

---

*Supported in part by the National Science Foundation.

[1]F. Bertaut and F. Forrat, Compt. rend. 242, 382 (1956).

[2]C. Herring and C. Kittel, Phys. Rev. 81, 869 (1951).

[3]W. P. Wolf and G. P. Rodrigue, J. Appl. Phys. 29, 105 (1958).

[4]J. S. Kouvel, Phys. Rev. 102, 1489 (1956).

[5]H. Suhl, Proc. Inst. Radio Engrs. 44, 1270 (1956); J. Chem. Phys. Solids 1, 1 (1957).

[6]C. Kittel, Phys. Rev. 110, 836 (1958).

[7]Akhiezer, Bar'iakhtar, and Peletminskii, J. Exptl. Theoret. Phys. U.S.S.R. 35, 228 (1958)[translation: Soviet Phys. JETP 8, 157 (1959)].

[8]H. J. McSkimin (private communication).

SUPERCONDUCTIVITY VERSUS FERROMAGNETISM IN LANTHANUM-GADOLINIUM ALLOYS

R. A. Hein and R. L. Falge, Jr.
United States Naval Research Laboratory, Washington, D. C.

and

B. T. Matthias and C. Corenzwit
Bell Telephone Laboratories, Murray Hill, New Jersey

(Received May 25, 1959)

The differential paramagnetic susceptibility of samples of LaGd alloys containing from 0.60 to 5.00% Gd has been measured in the temperature range from 4.22 to 0.10°K. Matthias et al.[1] have shown that the addition of small amounts of Gd to La produces solid solutions with unusual magnetic properties. Alloys containing up to 0.90% Gd are superconductors with transition temperatures (T_C) which decrease linearly with increasing Gd content. Alloys containing 3.0 or more percent Gd are ferromagnetic with Curie temperatures which increase with increasing Gd content. The alloy containing 3.0% Gd becomes ferromagnetic at 1.3°K, the lowest temperature they could attain. Extrapolation of these data suggests that the T_C vs percent Gd curve and the Curie temperature vs percent Gd curve would intersect for an alloy of 1.25% at a temperature of 0.5°K. Thus two possibilities are suggested: (1) the curves meet and terminate, indicating the possibility of a ferromagnetic superconductor for one specific composition, or (2) they intersect in such a way that there would exist a range of compositions for which a given alloy would exhibit (in zero magnetic field) both superconductivity and ferromagnetism, but at different temperatures. The primary purpose of the work reported here was to investigate the validity of these extrapolations.

Temperatures below 1°K were produced by the magnetic cooling method. The sample to be investigated was cemented in a copper holder which was in thermal contact with the paramagnetic salt via a copper rod 12 cm long and 3 mm in diameter. A carbon resistor cemented directly on the sample indicated that it was in intimate

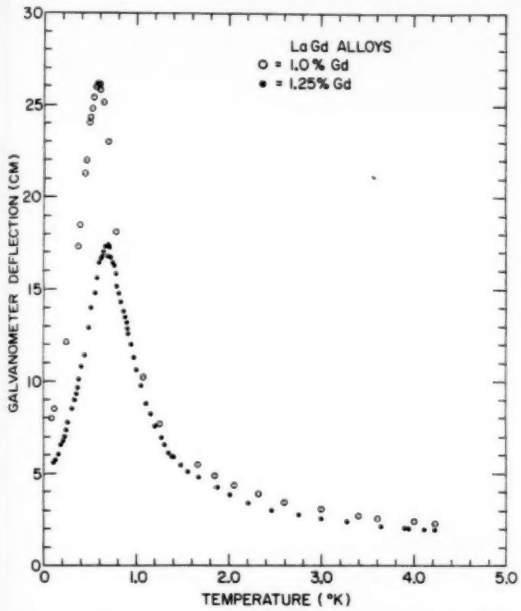

FIG. 1. Galvanometer deflection as a function of temperature obtained in zero applied magnetic field for the alloys containing 1.25 and 1.00% Gd.

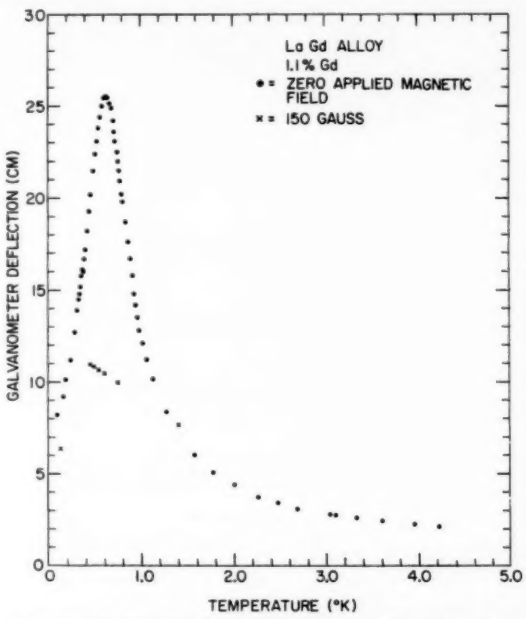

FIG. 2. Galvanometer deflection as a function of temperature and magnetic field for the alloy containing 1.10% Gd.

thermal contact with the salt.[2] The differential magnetic susceptibilities of the salt and sample were independently observed by means of a dc mutual inductance technique described in detail elsewhere.[3]

In Fig. 1 galvanometer deflections obtained in zero applied magnetic field with the 1.25 and 1.00% Gd alloys are plotted as a function of temperature. Such a plot represents susceptibility in arbitrary units vs temperature. These curves suggest the occurrence of a magnetic transition of some kind.

One possible interpretation of the curves in Fig. 1 is that the decrease in the susceptibility at the lower temperatures is due to superconductivity in parts of the sample. The entire sample could not be superconducting as the susceptibility was still paramagnetic at the lowest temperatures. If the decrease in susceptibility were caused by the diamagnetism of superconducting inclusions, the application of a magnetic field less than critical would have no effect on the susceptibility while a field greater than critical would cause an increase in the susceptibility. Application of a magnetic field actually caused the susceptibility to decrease, as shown in Fig. 2. Therefore, these alloys (1.00, 1.10, and 1.25% Gd) do not become superconducting in the earth's magnetic field down to a temperature of 0.1°K.[4]

The maxima in the curves of Figs. 1 and 2 can be interpreted as evidence of antiferromagnetic or ferromagnetic transitions. Although pure Gd is ferromagnetic, Thoburn et al.,[5] on the basis of magnetic moment measurements at low temperatures, reported antiferromagnetic transitions in LaGd alloys containing 47 to 90% Gd. On the other hand, Matthias et al.[1] observed large remanent moments for LaGd alloys containing 3 to 10% Gd, and concluded that these alloys are ferromagnetic below 6.0°K.[1] In keeping with this latter conclusion the temperatures of the maxima in the susceptibility have been taken, in this report, to be approximations to ferromagnetic Curie temperatures.[6] Curie temperatures so defined are plotted as a function of Gd content in Fig. 3.

The 0.60% Gd alloy becomes superconducting at 2.6°K.[1] Extrapolation of the Curie temperature curve in Fig. 3 indicates that this alloy might also become ferromagnetic at 0.3°K. The susceptibility of a sample of this alloy was paramagnetic at 4.2°K and increased as the temperature was lowered. A superconducting transition was observed which, in zero applied magnetic

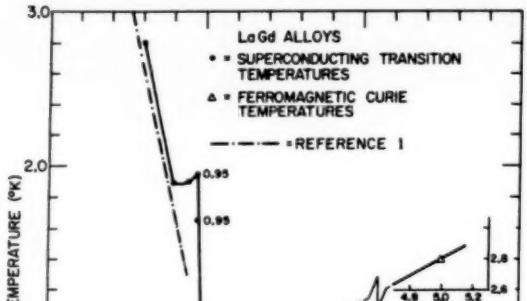

FIG. 3. Superconducting transition temperatures and ferromagnetic Curie temperatures of the La-Gd alloys as a function of Gd concentration.

field, was spread over the temperature interval from 2.8°K to 1.3°K. From 1.3°K down to 0.1°K the sample remained diamagnetic with a constant susceptibility. No evidence of ferromagnetism was observed in zero magnetic field.

Alloys in the range of 0.80 to 0.98% Gd were also studied. All samples were found to be superconducting. The zero field transitions (T_c) for the 0.60 and 0.80% Gd samples were in accord with the linear behavior previously reported.[1] However, departure from linear behavior was noted for a 0.95% sample. Its T_c was higher than expected. The first reaction to this result was to question the Gd content of the sample. Further doubt about composition arose when a 0.98% sample was found to possess a T_c which agreed with the linear dependence. To resolve this question of composition, a second 0.95% sample was fabricated. Its T_c was higher than that of the first. In fact, its T_c was even higher than that of the 0.80% sample. A 0.90% sample was then investigated and the results substantiated the deviation from linearity observed for the 0.95% samples. The dependence of T_c on Gd concentration is depicted in Fig. 3. In drawing the curve in this figure, more weight was given to the second 0.95% sample as extra care was taken in its fabrication.

In Fig. 3 one notes rather strange behavior in the vicinity of 0.95% Gd. Increasing the Gd content from 0.95 to 0.98% causes T_c to decrease by a factor of two. Increasing the Gd content from 0.98 to 1.00% causes the sample to exhibit ferromagnetism rather than superconductivity. On the other hand, decreasing the Gd content from 0.95 to 0.80% has relatively little effect on T_c. The sharp decrease in T_c observed for those alloys between 0.95 and 0.98% is consistent with the ideas of Suhl and Matthias.[7] They predicted that the slope of T_c vs percent Gd would approach infinity as T_c approaches zero. The observed slope is 30°K per percent Gd. The behavior of T_c in the region between 0.80 and 0.95% and the change from superconductivity to ferromagnetism between 0.98 and 1.00% are not yet understood.

Measurements of the susceptibility of the 0.95% sample in magnetic fields sufficient to destroy superconductivity suggest that this sample is ferromagnetic at the fields and temperatures involved. Critical magnetic field measurements, as well as magnetic moment measurements, will be included in a more comprehensive report.

---

[1]Matthias, Suhl, and Corenzwit, Phys. Rev. Letters 1, 92 (1958).

[2]R. A. Hein, Phys. Rev. 102, 1511 (1956).

[3]D. DeKlerk, Handbuch der Physik (Springer-Verlag, Berlin, 1956), Vol. 15, p. 74.

[4]Since this work was initiated an overlap between superconductivity and ferromagnetism has been reported by Matthias, Suhl, and Corenzwit, Phys. Rev. Letters 1, 449 (1958).

[5]Thoburn, Legvold, and Spedding, Phys. Rev. 110, 1298 (1958).

[6]R. M. Bozorth, Ferromagnetism (D. Van Nostrand Company, Inc., New York, 1951), p. 56.

[7]H. Suhl and B. T. Matthias, Phys. Rev. Letters 2, 5 (1959).

LONGITUDINAL MAGNETORESISTANCE OF METALS IN HIGH FIELDS.

B. Lüthi
Institut für Kalorische Apparate und Kältetechnik, Swiss Federal Institute of Technology, Zürich, Switzerland
(Received May 4, 1959)

Recently de Launay, Dolecek, and Webber[1] have measured the longitudinal resistance change of polycrystalline copper in magnetic fields up to 100 kilogauss, and they found a saturation although the transverse magnetoresistance appears to increase linearly without limit. We want to present, here, preliminary results of longitudinal magnetoresistance measurements in high magnetic fields on a number of other polycrystalline metals which show a similar effect.

Magnetic fields up to 200 kilogauss were produced by condenser discharges through coils cooled by liquid air or liquid helium.[2] Pulse lengths between 5 msec and 30 msec were available.[3] A number of experimental problems must be solved for such measurements. Shaking of the specimen due to shaking of the field coil and due to the Lorentz forces was prevented by mounting the specimen very carefully; the eddy-current heating was diminished by taking wires with only 0.2-mm diameter, and by choosing an appropriate discharge time. Furthermore, care had to be taken to place the specimen parallel to the external field, since the transverse magnetoresistance can be orders of magnitude bigger than the longitudinal one. The ratios of the transverse to the longitudinal magnetoresistance in the fields attainable with our apparatus are, for example, for Li, Ag, Au, and Pt about 10, for Sn 10^2, and for Pb and Zn about 10^3. As a control of our technique, the results on some metals were checked by means of dc measurements in low fields on specimens of extreme purity.

Figure 1 shows the results of Li, Ag, Au, Zn, Sn, Ni, Pt, and Pb in a Justi-Kohler plot, in which ΔR means the resistance change, R_T the zero-field resistance at temperature T (liquid air or liquid helium) and R_θ the resistance at the Debye temperature. The results plotted in this figure show quite clearly that there exists a very strong trend to saturation in the high-field region for all metals measured so far except Li, which shows a strict linear behavior, and Pb, which follows an $H^{1/2}$ law. Bearing in mind the above-mentioned experimental difficulties, one may state that the longitudinal magnetoresistance of a great number of polycrystalline metals shows a very weak field dependence in high magnetic fields.

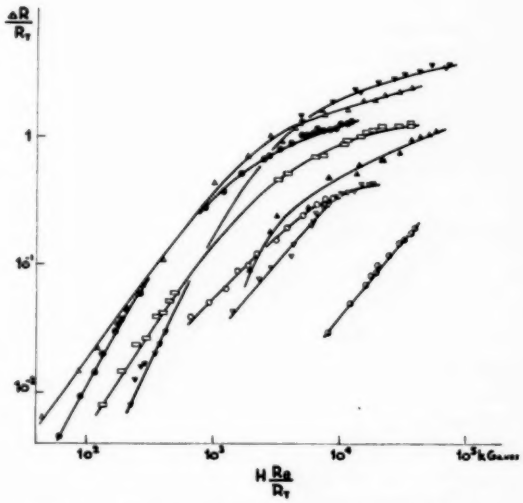

FIG. 1. Longitudinal magnetoresistance of

◕ Li (430) ▭ Ag (220) ● Au (320)
△ Zn (240) ▼ Sn (160) ○ Ni (410)
▽ Pt (240) ▲ Pb (90).

The figures in brackets indicate the Debye θ values used in calculating R_θ.

For single crystals this behavior would be in nice agreement with existing theories.[4] However, as Ziman[5] has pointed out, in the case of polycrystalline metals one has to take an average over all microcrystal directions, and here the special form of the energy surface plays an important role. Ziman's model with some very special assumptions gives a linear dependence for the transverse and the longitudinal magnetoresistance, but this model with its cylindrical surface is too special to account for a real metal such as Li. For the other monovalent metals Cu, Ag, and Au, one can give an explanation of the magnetoresistance behavior (linear dependence for the transverse case, saturation for the longitudinal one) by allowing the Fermi surface to touch the Brillouin-zone face. However, since our results show that a saturation-like behavior of the longitudinal magnetoresistance appears to occur for a wide range of metals and with only few exceptions, an explanation inde-

pendent of the detailed form of the energy surfaces is needed. These measurements are being continued on other substances.

I wish to thank Professor P. Grassmann, the Director of this Institute, and Dr. J. L. Olsen for their kind encouragement and support. This work was supported financially by a Arbeitsbeschaffungskredit des Bundes.

[1]de Launay, Dolecek, and Webber, J. Phys. Chem. Solids (to be published).

[2]J. L. Olsen, Helv. Phys. Acta 26, 798 (1953).

[3]The measurements with longer pulse length were carried out using apparatus kindly lent by Mr. P. Cotti.

[4]Lifshits, Azbel', and Kaganov, J. Exptl. Theoret. Phys. U.S.S.R. 31, 63 (1956) [translation: Soviet Phys. JETP 4, 41 (1957)].

[5]J. M. Ziman, Phil. Mag. 3, 1117 (1958).

STIMULATED EMISSION OF RADIATION BY RELATIVISTIC ELECTRONS IN A MAGNETIC FIELD*

Jürgen Schneider
Department of Physics, Duke University, Durham, North Carolina
(Received April 27, 1959)

The Schrödinger equation of an electron, moving perpendicularly to a magnetic field H, can be reduced to the equation of the harmonic oscillator.[1] The kinetic energy W_i of an electron in the state i is therefore given by $W_i=(i+\frac{1}{2})\hbar\omega_0$, where $\omega_0=eH/m_0c$ is the cyclotron frequency, associated with all transitions between the states i and $i+1$. Using the relativistic Schrödinger equation and neglecting the spin, the kinetic energy levels of the electron are found to be[2]

$$W_i=m_0c^2[1+2(i+\tfrac{1}{2})\hbar\omega_0/m_0c^2]^{\frac{1}{2}}-m_0c^2. \quad (1)$$

Transitions between states $i+1$ and i now result in a whole spectrum of frequencies $\omega_{i+1,i}$:

$$\omega_{i,i+1}=\omega_0(1-i\hbar\omega_0/m_0c^2). \quad (2)$$

Second order terms in $\hbar\omega_0/m_0c^2$ have been neglected. Consider N_i electrons in the energy state i. Under the influence of an alternating electric field they will undergo absorption and induced emission at the frequencies $\omega_{i,i+1}$ and $\omega_{i,i-1}$. The net transfer of power is

$$P_i=N_i\hbar(\omega_{i,i+1}w_{i,i+1}-\omega_{i,i-1}w_{i,i-1}). \quad (3)$$

The transition probability $w_{i,i+1}=w_{i+1,i}$ is given by

$$w_{i,i+1}=\hbar^{-2}E^2(\mu_{i,i+1})^2g(\omega). \quad (4)$$

The quantity

$$\mu_{i,i+1}=e[(i+1)\hbar/2m_0\omega_0]^{\frac{1}{2}} \quad (5)$$

is the matrix element of the harmonic oscillator for the transition between the state i and $i+1$; E is the amplitude and ω the frequency of the external electric field. For sufficiently sharp lines, i.e., for a small collision frequency $1/\tau$, the response curve $g(\omega)$ can be assumed to be Lorentzian:

$$g(\omega)=\tau/[1+(\omega_{i,i+1}-\omega)^2\tau^2]. \quad (6)$$

Introducing Eqs. (2), (4), (5), and (6) into Eq. (3), one obtains

$$P_i=N_i(e^2\tau/m_0)(E^2/2)\varphi,$$

$$\varphi=\frac{(1-i\alpha)(i+1)}{1+(\omega_{i,i+1}-\omega)^2\tau^2}-\frac{[1-(i-1)\alpha]i}{1+(\omega_{i,i-1}-\omega)^2\tau^2}, \quad (7)$$

where $\alpha=\hbar\omega_0/m_0c^2$. Neglecting again second order terms in α, the function φ can be written

$$\varphi(x)=\frac{1}{1+x^2}+2Q\frac{W}{m_0c^2}\frac{x}{(1+x^2)^2}, \quad (8)$$

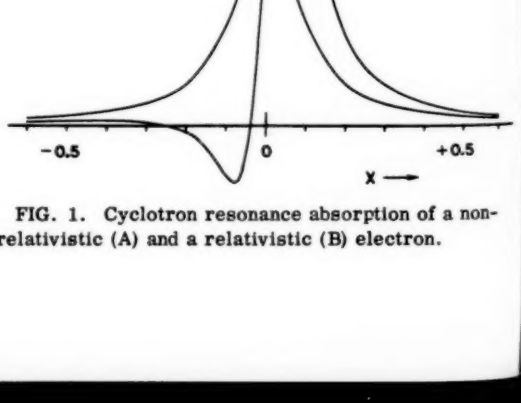

FIG. 1. Cyclotron resonance absorption of a nonrelativistic (A) and a relativistic (B) electron.

where $x=(\omega_{i,i+1}-\omega)\tau$, W is the kinetic energy of the electron, and $Q=\omega_0\tau$. In the nonrelativistic limit, i.e., for $W/m_0c^2=0$, the cyclotron resonance line is Lorentzian. However, for $QW/m_0c^2>1$, the function φ may become negative, corresponding to a net stimulated emission instead of absorption. This effect is shown in Fig. 1 where we have taken $QW/m_0c^2=1.5$, e.g., $Q=10\,000$ and $W=76$ ev.

If the electrons are not monochromatic, but have rather a distribution of energies, one obtains essentially the same effect, if there is an overpopulation of the upper states.[2] The above formulas (7) and (8) have also been derived[2] by means of the Boltzmann transfer equation, taking into account the dependence of the electron mass on the kinetic energy.

It does not appear unlikely that this effect could be used for a new type of maser, which would require no microwave "pump" and no low-temperature operation.[2]

The author wishes to thank Dr. Walter Gordy for his interest in this work.

*This research was supported by the U. S. Air Force through the Air Force Office of Scientific Research of the Air Research and Development Command.

[1]L. Landau, Z. Physik 64, 629 (1930).

[2]J. Schneider (to be published).

STRENGTHS OF HIGH-ENERGY CAPTURE GAMMA RAYS IN W^{184}

D. J. Hughes, M. K. Brussel, J. D. Fox, and R. L. Zimmerman
Brookhaven National Laboratory, Upton, New York
(Received May 27, 1959)

During the past few years the Brookhaven fast chopper[1] has been devoted to an extensive study of the parameters of energy levels in heavy nuclides excited by the capture of slow neutrons. The results, obtained primarily from total cross sections alone, have given information on the distribution laws of radiation widths, neutron widths, and spacings of these energy levels.[2] However, little is known of the spins, J, of the levels, aside from the limitation to two values, $I\pm\frac{1}{2}$, with I the target nucleus spin. In a few favorable cases, for which Γ_n is larger than Γ_γ, the total cross section suffices for J determination, but in general, partial cross sections, scattering or capture, must be measured as well.

Another possible method for spin measurement is based on the capture gamma rays emitted from the individual levels. The most direct use of capture gamma rays for determination of J applies when the ground-state transition is allowed for one of the spins only.[3] An example[4] is the compound nucleus W^{184}, in which the electric dipole ground-state transition, of energy[5] 7.42 Mev, is allowed for the $J=1$ but forbidden for the $J=0$ levels. In this ground-state technique of spin identification the question of the variation in transition probability from level to level is important, for it might be possible that a $J=1$ level would exhibit no observable ground-state gamma because of an extremely weak transition probability. Because of this application to spin determination, and the intrinsic interest in transition probabilities as well, it is desirable to gain information on the high-energy capture gamma rays in heavy nuclides.

Essentially no information exists concerning the relative transition probabilities involved in the present situation, that is, the emission of gamma rays from neighboring levels a few ev apart to the same final state, some 6 or 7 Mev distant. The typical neutron capture gamma-ray measurements do not supply the necessary information, for they are limited to excitation by one energy only, usually thermal neutrons. Electric dipole transitions in the hundred-kev range[6] show wide variations in transition probability and are usually hindered by large amounts relative to the high-energy capture gamma rays; the latter are in reasonable agreement with estimates[6] based on Weisskopf's single-particle model, although exhibiting variations of the order of a factor of ten from nuclide to nuclide.

Earlier capture gamma measurements with the Brookhaven fast chopper showed[4] strong ground-state transitions[7] for several of the known resonances in the target nucleus W^{183}, indicating that $J=1$ for these resonances. In that work gamma-ray energies were measured in sodium iodide crystals located near the capturing sample, and neutron energies simultaneously by time of flight to the sample. The same equipment has now been installed at the NRU reactor, Chalk River, Can-

ada, where the fast chopper flux is some 80 times greater and correspondingly higher resolution is possible. Definite ground-state gamma rays have now been detected in five of the six known levels up to an energy of 100 ev; thus these levels are identified as $J=1$. The spins are in agreement with the results of Waters et al.[8] based on partial cross sections, except for the 41- and 66-ev levels, too weak to be analyzed by partial cross-section methods. It is noteworthy that ground-state gammas can be detected by the present method for such weak levels.

The transition probabilities of the ground-state gamma rays for the five $J=1$ levels are given in Table I relative to the average value, taken as unity. The absolute value of this average probability, obtained by comparison with thermal neutron results,[5,9] is 2.6×10^{-3} ev, expressed as a partial radiation width. The striking fact about the transition probabilities is their small variation from level to level, an average deviation from the mean of only 20% which is very different from the wide variation of neutron widths, also listed in Table I for comparison purposes. These results are surprising in the light of the discussion of Porter and Thomas,[10] who assume that the individual gamma rays would show an extreme variation from level to level of the same type as that observed[11] for neutron widths. Although the number of levels in the present work is small as yet, the close grouping of the observed transition probabilities is suggestive that the high-energy radiation from neighboring levels in heavy nuclides does not exhibit a range as wide as the Porter-Thomas neutron width distribution. Measurements for other similar nuclides are continuing because of the intrinsic interest of the transition probabilities and for the determination of spins of more levels.[12] We wish to acknowledge the cooperation of the Chalk River scientists in setting up the cooperative program that made these results possible, particularly H. B. Lewis and C. H. Westcott. H. Palevsky and R. E. Chrien have given valuable help in the design of the equipment and W. C. Olsen in its operation.

Table I. Relative transition probabilities of ground-state gamma rays emitted from specific $J=1$ levels in W^{184}. The levels are identified by the energy of the corresponding neutron resonance and the reduced neutron widths of these resonances are also listed.

Neutron energy (ev)	Reduced neutron widths (10^{-3} ev)	Relative transition probability (mean value = 1.00)
7.6	1.2	1.27 ± 0.24
27.1	9.2	1.01 ± 0.03
41	0.16	0.53 ± 0.12
47	37	1.00 ± 0.04
66	0.26	1.19 ± 0.27

[1]Seidl, Hughes, Palevsky, Levin, Kato, and Sjöstrand, Phys. Rev. 95, 476 (1954).

[2]For a summary see D. J. Hughes, Neutron Cross Sections (Pergamon Press, New York, 1957).

[3]H. H. Landon and R. Rae, Phys. Rev. 107, 1333 (1957).

[4]Fox, Zimmerman, Hughes, Palevsky, Brussel, and Chrien, Phys. Rev. 110, 1472 (1958).

[5]B. B. Kinsey and G. A. Bartholomew, Can. J. Phys. 31, 1051 (1953).

[6]B. B. Kinsey, Handbuch der Physik (Springer-Verlag, Berlin, 1957), Vol. 40.

[7]These might include the similar allowed transitions to the 2^+ state at 111 kev, but none of the present conclusions would be significantly affected.

[8]Waters, Evans, Kinsey, and Williams, Nuclear Phys. (to be published).

[9]Groshev, Demidov, Lutsenko, and Plevhanov, Atlas of Capture Gamma Rays (Moscow, 1958).

[10]C. E. Porter and R. G. Thomas, Phys. Rev. 104, 483 (1956).

[11]D. J. Hughes and J. A. Harvey, Phys. Rev. 99, 1032 (1955).

[12]Similar measurements of ground-state transitions for the same five levels in W^{184} have recently been reported by Huynk, Julien, Corge, Netter, and Simic, Compt. rend. 248, 2330 (1959). Their results and those of Table I are in excellent agreement. In addition, J. R. Bird has reported qualitative results for these levels in the Proceedings of the Second United Nations International Conference on the Peaceful Uses of Atomic Energy, Geneva, 1958 (United Nations, Geneva, 1959), paper P/35.

SYSTEM OF LEVELS IN EVEN-EVEN NUCLEI*

C. A. Mallmann
Argonne National Laboratory, Lemont, Illinois
(Received May 11, 1959)

It is of interest, in order to stimulate further experimental and theoretical work, to point out the existence of a system of levels in even-even nuclei whose energies are related by an empirically determined law which seems to be applicable throughout the nuclear chart.

If one plots the ratio R of the energy of the first excited 6+ (denoted ^{1}6+) and 8+ (^{1}8+) states to the energy of the first excited 2+ (^{1}2+) state as a function of the ratio $R(^1\text{4+})$ of the energy of the first excited 4+ (^{1}4+) state to the energy of the first excited 2+ (^{1}2+) state, one finds that practically all the points lie on two curves, one for the ^{1}6+ states and one (tentative) for the ^{1}8+ states; see Fig. 1. Some of the energy ratio [for $R(^1\text{4+})$ larger than 2.9] were taken from Scheuer and Aisenberg's paper[1] with the exception of Gd^{158}.[2] All the values used are given in Table I. Question marks in Fig. 1 mean that there is uncertainty in the assignment of the levels considered. We also show in Fig. 1 the obvious lines for $R(^1\text{4+})$ and $R(^1\text{2+})=1$ in order to visualize better the position of the ^{1}2+, ^{1}4+, ^{1}6+, ^{1}8+ levels. Naturally the ^{1}0+ ground state is at $R=0$.

As one would expect, the $R(^1\text{8+})$ and $R(^1\text{6+})$ curves coincide over a range of $R(^1\text{4+})$ values with the straight lines $R(^1\text{8+})=-312/7+(594/35)R(^1\text{4+})$ and $R(^1\text{6+})=-11+(27/5)R(^1\text{4+})$. These are a direct consequence of

$$E_I = aI(I+1) - b[I(I+1)]^2, \qquad (1)$$

which is the expression for the energy of a rotational state in an axially symmetric nucleus[3] taking into account the rotational-vibrational[4] interaction, where a and b are constants and I the spin of the state. The coefficients in the expressions for $R(^1\text{4+})$ and $R(^1\text{6+})$ are independent of a and b. The curves show that the rotational nuclei are confined to $R(^1\text{4+})$ values larger than or equal to 3.27 because for values smaller than this the empirical curves deviate from (1). Of the rotational nuclei, Hf^{178}, U^{232}, U^{236}, and Pu^{240} do not fall on the theoretical curve (1) within the experimental errors.

The theory of Davydov and Filippov[5] which treats rotations of nonaxially-symmetric nuclei, extends the agreement between theoretical and empirical curves down to values of $R(^1\text{4+})=8/3$ because $R(^1\text{4+})$ and $R(^1\text{6+})$ are reduced from 10/3 and 7, respectively, for $\gamma=0°$ to 8/3 and ~5.0 for $\gamma=30°$. Unfortunately the precise predictions for $R(^1\text{6+})$ have not been published so that a more careful comparison with experimental values cannot be made. The $R(^2\text{2+})$ (second 2+ excited state) values

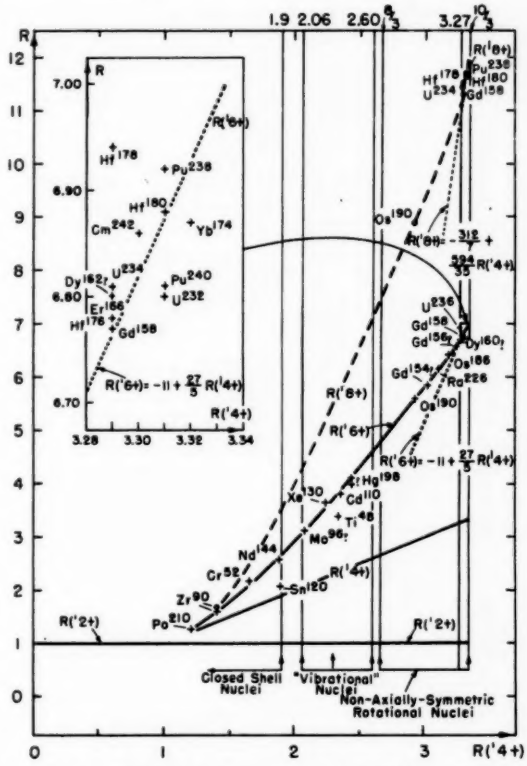

FIG. 1. The figure shows the ratios $R(^1\text{8+})$, $R(^1\text{6+})$, $R(^1\text{4+})$, and $R(^1\text{2+})$ as a function of the ratio $R(^1\text{4+})$. Experimental points are indicated with a circle for $R(^1\text{8+})$ and a plus sign for $R(^1\text{6+})$. The $R(^1\text{6+})$ points for $R(^1\text{4+})$ values in the range from 3.28 to 3.33 are shown in a magnified scale at the upper-left side of the figure. Experimental errors are not indicated in order to avoid confusion in the graph. They are given in Table I. Question marks mean that there is an uncertainty in the assignment of the levels considered. Full lines are experimental; the dashed line is tentative for $R(^1\text{8+})$. The dotted lines are the theoretical predictions for rotational states in axially symmetric nuclei taking into account the rotational-vibrational interaction.

Table I. Experimental energy ratios with their errors used in Fig. 1.

	$E(^12+-{}^10+)$ (Mev)	$R(^14+)$	$R(^16+)$	$R(^18+)$
Po^{210}	1.183	1.21 ± 0.01	1.25 ± 0.01	...
Zr^{90}	2.197	1.41 ± 0.03	1.57 ± 0.03	1.64 ± 0.04
Cr^{52}	1.433	1.65 ± 0.01	2.17 ± 0.02	...
Sn^{120}	1.170	1.89 ± 0.02	2.06 ± 0.02	...
Nd^{144}	0.696	1.89 ± 0.01	2.56 ± 0.02	...
Mo^{96}	0.778	2.08 ± 0.03	3.10 ± 0.04	...
Xe^{130}	0.534	2.24 ± 0.02	3.64 ± 0.03	...
Ti^{48}	0.986	2.33 ± 0.02	3.39 ± 0.02	...
Cd^{110}	0.656	2.35 ± 0.02	3.78 ± 0.03	...
Hg^{198}	0.41176	2.43 ± 0.02	3.97 ± 0.03 } ? 4.09 ± 0.03 }	...
Gd^{154}	0.1229	3.020 ± 0.005	5.83 ± 0.01	...
Gd^{156}	0.08897	3.240 ± 0.005	6.57 ± 0.01	...
Gd^{158}	0.07956	3.292 ± 0.003	6.774 ± 0.006	11.30 ± 0.02
Dy^{160}	0.087	3.27 ± 0.01	6.69 ± 0.01	...
Dy^{162}	0.0808	3.290 ± 0.005	6.80 ± 0.02	...
Er^{166}	0.0805	3.290 ± 0.005	6.80 ± 0.08	...
Yb^{174}	0.0766	3.32 ± 0.02	6.87 ± 0.03	...
Hf^{176}	0.08835	3.290 ± 0.005	6.77 ± 0.13	...
Hf^{178}	0.09319	3.29 ± 0.01	6.94 ± 0.03	11.50 ± 0.09
Hf^{180}	0.09328	3.31 ± 0.01	6.88 ± 0.02	11.64 ± 0.06
Os^{186}	0.13721	3.18 ± 0.06	6.39 ± 0.12	...
Os^{190}	0.187	2.92 ± 0.02	5.59 ± 0.06	8.88 ± 0.08
Ra^{226}	0.06776	3.10 ± 0.01	6.14 ± 0.02	...
U^{232}	0.0472	3.31 ± 0.02	6.80 ± 0.04	...
U^{234}	0.04350	3.29 ± 0.02	6.81 ± 0.03	11.47 ± 0.09
U^{236}	0.04528	3.26 ± 0.03	6.92 ± 0.12	...
Pu^{238}	0.04407	3.310 ± 0.005	6.92 ± 0.02	11.68 ± 0.23
Pu^{240}	0.04287	3.310 ± 0.005	6.810 ± 0.005	...
Cm^{242}	0.04212	3.30 ± 0.01	6.86 ± 0.12	...

of Davydov and Filippov as a function of $R(^14+)$ values seem to be in agreement with experimental results. The agreement of the theoretical predictions for transition probability ratios and $E2$, $M1$ admixtures in the $^22+-{}^12+$ transition is quite good.[6,7] It is interesting to note that the smallest $R(^14+)$ predicted by this theory is 8/3; clearly many nuclei have $R(^14+)$ smaller than 8/3 (see Fig. 1) so that they are not in agreement with this theory.

The "vibrational" nuclei[8] have average $R(^14+)$ value of 2.3[9] with values ranging from 2.07 to 2.63 and the closed shell nuclei have $R(^14+)$ values smaller than 1.9.[9] The nuclei above the "vibrational" nuclei have been called nonaxially-symmetric rotational nuclei in Fig. 1. Of all the nuclei which are not nonaxially-symmetric rotational nuclei only Ti^{48} and Sn^{120} do not fall within their experimental error on the empirical curve.

The classification of nuclei into different types seems not to be essential, at least with respect to this system of levels, because all nuclei follow the same curves which relate the energies of these excited states.

The author wishes to thank F. T. Porter and M. S. Freedman for a critical reading of the manuscript.

---

*This work was performed under the auspices of the U. S. Atomic Energy Commission.

[1]W. Scheuer and E. Aisenberg, Atomic Energy Commission, Argentina, 1958 (to be published).

[2]Knowles, Bartholomew, and Campion, Bull. Am. Phys. Soc. Ser. II, 4, 246 (1959).

[3]A. Bohr and B. R. Mottelson, Phys. Rev. 90, 717 (1953).

[4]Alder, Bohr, Huus, Mottelson, and Winther, Revs. Modern Phys. 28, 432 (1956).

[5]A. S. Davydov and G. F. Filippov, Nuclear Phys. 8, 237 (1958).

[6]D. M. Van Patter, Bull. Am. Phys. Soc. Ser. II, 4, 233 (1959).
[7]Potnis, Malik, and Mandeville, Bull. Am. Phys. Soc. Ser. II, 4, 233 (1959).
[8]G. Scharff-Goldhaber and J. Weneser, Phys. Rev. 98, 212 (1955).
[9]C. A. Mallmann, Proceedings of the Second United Nations International Conference on the Peaceful Uses of Atomic Energy, Geneva, 1958 (United Nations, Geneva, 1959), paper A/Conf. P/1971.

ANTIPARTICLES IN SPACE

Edward W. Hones, Jr.
Physics Section, Convair,* San Diego, California
(Received May 11, 1959)

About a decade ago, Hoyle[1] obtained a steady-state model for an expanding universe by introducing continuous creation of matter into the field equations of general relativity. More recently, Morrison[2] has speculated that the creation of matter throughout the universe may take place by formation of particle, antiparticle pairs which subsequently are segregated by a repulsive force between matter and antimatter. Burbidge and Hoyle[3] have estimated that several observed astrophysical phenomena (both within our galaxy and outside it) might be explained as due to the presence of antimatter. But they conclude that the ratio of antimatter to ordinary matter in our galaxy cannot exceed $\sim 10^{-7}$. The purposes of this Letter are (a) to point out that a slightly more stringent limit for this ratio is suggested by cosmic-ray measurements reported by Perlow and Kissinger[4,5] and (b) to observe that relatively simple equipment in earth satellites or in space probes should serve either to detect antiparticles in our galaxy or to define a much more stringent limit on their number density than can presently be set.

If one assumes that antiprotons exist throughout the galaxy in the limiting number density, n_-, of $10^{-7}/\text{cm}^3$ estimated by Burbidge and Hoyle and if one takes the $p,\bar{p}$ annihilation cross section used by them and adopts, for protons in the galaxy, the number density, n_+, of 1 per cm^3, one obtains a $p,\bar{p}$ annihilation rate of 3×10^{-22} cm^{-3} sec^{-1}. From an average annihilation, there ultimately appear 3.6 photons of ~170 Mev and 3.6 photons of 0.51 Mev.[6] Therefore, the source of gamma rays of each type has an intensity, throughout the galaxy, of 10.8×10^{-22} cm^{-3} sec^{-1}. The flux, Q (cm^{-2} sec^{-1}) of these photons coming from a spherical region of radius R (cm) surrounding the earth is

$$Q=\tfrac{1}{2}PR \text{ cm}^{-2} \text{ sec}^{-1}, \quad (1)$$

where P is the photon generation rate in the source (cm^{-3} sec^{-1}). We can obtain, simply, a very conservative value for the flux, by calculating those coming from a sphere of 10 000 light years (10^{22} cm) radius. (The flux from the whole galaxy would certainly be several times larger than this.) We find, for each type of photon (170 Mev and 0.51 Mev),

$$Q=\tfrac{1}{2}(10.8\times10^{-22}\times10^{22})=5.4 \text{ photons/cm}^2 \text{ sec}.$$

In terms of energy flux, this is ~900 Mev/cm^2 sec for the 170-Mev photons and ~2.5 Mev/cm^2 sec for the 0.51-Mev photons.

Perlow and Kissinger measured the primary gamma radiation at 80 km altitude and found the flux of photons in the energy range 0.1 to 15 Mev to be $\leqq 2.7/\text{cm}^2$ sec from the upper hemisphere.[4] This figure seems to be in fair agreement with the calculated flux of 0.51-Mev photons until one recalls that (a) the calculated flux is not that from the whole galaxy but only from a 10^{22}-cm sphere and (b) the measurements included not only 0.51-Mev photons but all photons between 0.1 and 15 Mev. From these considerations one is probably justified in saying that $n_-/n_+ < 10^{-8}$ in the galaxy. Perlow and Kissinger measured, also, the flux of gamma rays of energy between 3.4 and 90 Mev and found an energy flux of 1.4 Mev/cm^2 sec from the upper hemisphere.[5] It is not possible to make a very meaningful comparison of these data with the calculated flux of 170-Mev gamma rays because the latter (if they occur) would not have been detected very efficiently by the apparatus used. However, the fact that evidence of this large energy flux (~900 Mev/cm^2 sec) has not been found in other cosmic-ray experiments indicates that the flux of ~170-Mev gammas is at least one or two orders of magnitude smaller than that calculated above (5.4/cm^2 sec).

Considerably more precise information concerning the density of antiparticles in the galaxy can be obtained by carrying a simple scintillation gamma-ray spectrometer some distance from the earth in a satellite or space probe to look specifically for 0.51-Mev photons.[7] A search could be made, also, for the high-energy (~170 Mev) gamma rays, although the spectrometer to be used for these would be larger and more complicated.

An explicit search for positrons in space also can be made quite simply, using an experimental arrangement such as that shown schematically in Fig. 1. With this equipment one would detect positrons arriving at the annihilation block by the resulting coincident 0.51-Mev light pulses in the two NaI crystals. Using the interplanetary cosmic-ray flux found by Van Allen[8] and assuming a resolving time of 10^{-7} sec for NaI crystals, I estimate that at distances from the earth greater than ~20 000 miles this type of instrument could detect positrons if their number density were $10^{-14}/cm^3$ or greater.

Detection of positrons in space in the neighborhood of the earth's orbit would not necessarily be firm proof of the existence of antimatter on a cosmic (or even galactic) scale, because positrons from cosmic-ray showers or from other processes in the solar system might find their way to this region. Thus, the search for the characteristic annihilation gamma rays probably represents the better test for the widespread existence of antiparticles.

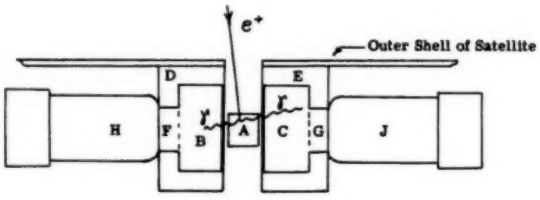

FIG. 1. Schematic representation of satellite instrumentation to detect positrons in space. A—annihilation block of low-Z material; B, C—NaI crystals shielded from sun by ~1 mg/cm² foil over surface; D, E—lead shields to reduce individual counting rates of the two NaI crystals; F, G—light pipes; H, J—photomultipliers. The system would include electronic circuits which would register only coincident 0.51-Mev pulses in crystals B and C.

*A division of General Dynamics Corporation.

[1]F. Hoyle, Monthly Notices Roy. Astron. Soc. 108, No. 5, 372 (1948).

[2]P. Morrison, Am. J. Phys. 26, 358 (1958).

[3]G. R. Burbidge and F. Hoyle, Nuovo cimento 4, 558 (1956).

[4]G. J. Perlow and C. W. Kissinger, Phys. Rev. 84, 572 (1951).

[5]G. J. Perlow and C. W. Kissinger, Phys. Rev. 81, 552 (1951).

[6]E. Segrè, Annual Review of Nuclear Science (Annual Reviews, Inc., Palo Alto, 1958), Vol. 8, p. 127.

[7]B. Rossi has informed me that his group is currently planning a satellite experiment to measure the gamma radiation in space.

[8]J. A. Van Allen and L. A. Frank, Nature 183, 430 (1959).

K-MESON NUCLEON INTERACTION*

Robert Karplus, Leroy Kerth, and Thaddeus Kycia
Department of Physics and Lawrence Radiation Laboratory,
University of California, Berkeley, California
(Received May 15, 1959)

It is to be expected that the experimental study of K-meson production, scattering, and absorption processes will shed light on the fundamental interactions of the strange particles. The most promising method of analysis for determining the coupling constants and relative parities of the particles exploits the analyticity properties of various reaction amplitudes.[1,2] The recently measured angular distribution of the K^+-proton scattering cross section at various intermediate energies[3] and the K^--proton scattering parameters[4,5] enable us to make a more accurate study of the K-nucleon-hyperon coupling terms than has previously been possible. Unfortunately, the accuracy is not sufficient to lead to a definite conclusion about the K-meson parity, but we do conclude that the K^--proton force is attractive. We hope that the present discussion also has value because it will make clear some of the difficulties that must be faced.

The numerical work is carried out most easily if one uses the observations to calculate the function $f(\omega)$ of K-meson laboratory energy ω, ($\hbar = c = m_K = 1$)

$$f(\omega) = \frac{1}{\pi}\int_1^\infty \frac{A_+(\omega')}{\omega' - \omega}d\omega' + \frac{1}{\pi}\int_{\omega_{\Lambda\pi}}^\infty \frac{A_-(\omega')}{\omega' + \omega}d\omega', \qquad (1)$$

where $A_\pm(\omega)$ are the imaginary parts of the $K^\pm$-proton forward scattering amplitudes and $\omega_{\Lambda\pi}$ is the (nonphysical) threshold for the reaction

$$K^- + p \to \Lambda + \pi. \qquad (2)$$

Since the integrals in Eq. (1) have been cut off, it is necessary to use a dispersion relation that does not weight the high-energy region. We have used this expression for the real parts $D_\pm(\omega)$ of the forward scattering amplitudes:

$$\omega_0 D_+(\omega) - \tfrac{1}{2}(\omega_0 + \omega)D_+(\omega_0) - \tfrac{1}{2}(\omega_0 - \omega)D_-(\omega_0)$$
$$\cong \omega_0 f(\omega) - \tfrac{1}{2}(\omega_0 + \omega)f(\omega_0) - \tfrac{1}{2}(\omega_0 - \omega)f(-\omega_0)$$
$$- \omega_0(\omega_0^2 - \omega^2)\left\{\frac{p_\Lambda X_\Lambda}{(\omega_\Lambda^2 - \omega_0^2)(\omega_\Lambda + \omega)} + \frac{p_\Sigma X_\Sigma}{(\omega_\Sigma^2 - \omega_0^2)(\omega_\Sigma + \omega)}\right\}. \qquad (3)$$

Here p_Λ (p_Σ) is the parity of $K^-\Lambda$ ($K^-\Sigma$) relative to the nucleon and X_Λ (X_Σ), proportional to the $K^-\Lambda N$ ($K^-\Sigma N$) coupling constant g_Λ^2 (g_Σ^2), is the magnitude of the residue at the pole $\omega = \omega_\Lambda$ (ω_Σ); these quantities are given by the expressions ($\lambda = \Lambda, \Sigma$)

$$\omega_\lambda = (m_\lambda^2 - m_p^2 - m_K^2)/2m_p,$$
$$X_\lambda = \frac{g_\lambda^2}{4\pi}\left|\frac{(m_\lambda + p_\lambda m_p)^2 - m_K^2}{4m_p m_\lambda}\right|, \qquad (4)$$

in terms of the masses m_λ, m_p, and m_K of the hyperons, proton, and K meson, respectively.

The energy ω_0 in Eq. (3) is one at which experimental information is available. By using a value substantially different from the meson rest energy, one can suppress the importance of the nonphysical region and of the amplitude $D_-(\omega_0)$, which decreases with energy.

The optical theorem is used to obtain $A_\pm(\omega)$ from the measurements of total cross section in the physical region. The K^+-proton total cross sections were taken from the data obtained by the use of nuclear emulsions[6] for energies below 200 Mev, the results of our experiment,[3] and the data of Burrowes et al.[7] for the higher energies. A curve was fitted to the data within an estimated statistical error of 10%. The low-energy K^--proton total cross sections have been measured in a hydrogen bubble chamber.[8] At higher energies the values are $\sigma_- = 52 \pm 9$ mb at $\omega = 2.08$,[9] and $\sigma_- = 60 \pm 20$ mb at $\omega = 2.52$.[10] We assumed the value $\sigma_- = 45$ at $\omega = 4$, and passed a smooth curve through all the points. The statistical error of the fit was estimated at 15%.

The values of $A_-(\omega)$ in the nonphysical region were obtained by extrapolating with the parameters of Dalitz and Tuan.[4,5] The integral is very sensitive to the sign of $D_-(1)$, the real part of the scattering amplitude at zero energy.[4] The combination of the integral and the scattering amplitude that occurs in Eq. (3), however, is not sensitive to this sign as long as the zero-range approximation gives the correct sign and a reasonably accurate value for $D_-(\omega_0)$. This is the case because the zero-range formula gives an analytic expression for the scattering amplitude which then essentially satisfies a dispersion relation by itself.

The results for the integration are given in Table I. They were obtained with a positive choice of $D_-(1)$. In the fourth column are the measured values[3] of $D_+(\omega)$; their magnitudes are simply related to the total cross sections, because it was found that the scattering is isotropic at $\omega = 1.46$ (225 Mev) and it was assumed that it is isotropic at the lower energies. In the last column is the value of D_- obtained from the zero-range parameters, which represent the data adequately.[4] An examination of the bubble chamber results[6] indicates that most of the elastic scattering is indeed a diffraction effect (i.e., the real part of the

Table I. Integrals and scattering amplitudes used in the dispersion relation Eq. (3). The K-meson energy is ω measured in units of the rest energy. The unit for the other quantities is the K-meson Compton wavelength.

ω	$f(\omega)$	$f(-\omega)$	$D_+(\omega)$	$D_-(\omega)$
1.00	3.0 ± 0.3	...	-1.25 ± 0.14	...
1.17	3.1 ± 0.3	5.2 ± 1.0	-1.24 ± 0.14	+0.56
1.285	3.1 ± 0.3	4.9 ± 0.8	-1.23 ± 0.14	+0.40
1.46	3.0 ± 0.3	...	-1.20 ± 0.08	...

scattering amplitude is small) and that there is accordingly some tendency toward forward scattering. Considering the experimental accuracy, one may estimate $|D_-(1.17)|<2$ and $|D_-(1.285)|<2$, consistent with Table I.

Since it is clearly not possible to solve for X_Λ and X_Σ separately, we have calculated the "average" coupling $\langle pX\rangle_{\mathrm{Av}}$ at an "average" pole $\bar{\omega}=\frac{1}{2}(\omega_\Lambda+\omega_\Sigma)$,

$$\langle pX\rangle_{\mathrm{Av}}=\tfrac{1}{2}(\bar{\omega}^2-\omega_0^2)(\bar{\omega}+\omega)\left[\frac{p_\Lambda X_\Lambda}{(\omega_\Lambda^2-\omega_0^2)(\omega_\Lambda+\omega)}+\frac{p_\Sigma X_\Sigma}{(\omega_\Sigma^2-\omega_0^2)(\omega_\Sigma+\omega)}\right], \quad (5)$$

for various choices of ω and ω_0 in Eq. (3). The results are listed in Table II. We conclude that this determination of $\langle pX\rangle_{\mathrm{Av}}$ leads to the value

$$\langle pX\rangle_{\mathrm{Av}}=0.0\pm0.5, \quad (6)$$

where we have included some provision for the uncertainty in $D_-(\omega_0)$.

Our conclusion, then, is that the K^+-proton interaction, at the present accuracy, is not sufficiently sensitive to the coupling parameters X_Λ and X_Σ to permit their evaluation from the present data. This is not too surprising, because Eq. (3) depends only on the difference of amplitudes at closely spaced energies that are quite distant from the poles ω_Λ and ω_Σ. We must also point out an important weakness of this type of analysis: only measurements on the $K^\pm$-neutron system can make possible a separation of the $K\Lambda P$ and $K\Sigma P$ coupling terms by their different isotopic properties. The determination of the K-meson parity from the angular dependence of associated $K\Lambda$ and $K\Sigma$ photoproduction[2] does not suffer from this defect and may, therefore, lead to a definitive result sooner.[11]

A few words are in order about the amplitude $D_-(\omega)$.[12] If it deviates substantially from its zero-range value at the energies of interest, then Eq. (6) will be modified somewhat [the right-hand side contains $D_-(\omega_0)$ with a coefficient -0.2]. The characteristic feature of the zero-range approximation, that the scattering is very strong in one of the two isotopic spin states,[4] depends essentially only on the fact that the large K^--p elastic cross section is accompanied by a substantial charge-exchange cross section at low energy. Each satisfactory set of parameters includes one scattering length whose real part exceeds 1.5 fermis in magnitude. We should like to add that only an extremely unlikely change in the experimental situation could reduce this value to 1 f. Since the elastic K^--proton interaction cannot have a range exceeding $\hbar/2m_\pi c$ ~0.7 f, it follows that the interaction is attractive in the isotopic state with the large scattering length. If $D_-(1)$ is negative, then the attraction is so strong as to give rise to a "bound" state which will be unstable to decay into a pion-hyperon system. $D_-(\omega)$ can be expected to change sign near $\omega\sim1.2$, a meson laboratory momentum of 300 Mev/c. If $D_-(1)$ is positive, then the attraction is not strong enough to give rise to a "bound" state. In this case $D_-(\omega)$ will decrease monotonically from its low-energy value; near $\omega\sim1.2$, it will already be quite small. Since the forces are attractive and strong in both of these cases, one cannot expect to distinguish between them by a study of the nuclear "optical model" potential for K^- mesons, which is then not simply related to the low-energy scattering amplitude. A measurement of the Coulomb interference effects in the angular distribution of K^--proton scattering will, of course, decide the matter.

Table II. "Average" couplings determined from Eq. (3) by the use of data at the energies ω and ω_0. The results are independent, to the accuracy of this work, of the choice of sign for $D_-(\omega)$.

ω	ω_0	$\langle pX\rangle_{\mathrm{Av}}$
1.46	1.285	-0.2 ± 0.6
1.46	1.17	-0.1 ± 0.4
1.00	1.285	$+0.1\pm0.4$
1.00	1.17	$+0.3\pm0.6$

*This work was carried out in part under the auspices of the U. S. Atomic Energy Commission and in part supported by the U. S. Air Force, Office of Scientific Research of the Air Research and Development Command.

[1]P. T. Matthews and A. Salam, Phys. Rev. 110, 565, 569 (1958); C. Goebel, Phys. Rev. 110, 572 (1958); K. Igi, Progr. Theoret. Phys. (Kyoto) 19, 238 (1958); John G. Taylor, Nuclear Phys. 9, 357 (1959).

[2]M. J. Moravcsik, Phys. Rev. Letters 2, 352 (1959).

[3]Kerth, Kycia, and Baender, Bull. Am. Phys. Soc. Ser. II, 4, 25 (1959); T. F. Kycia, thesis, University of California Report UCRL-8753 (unpublished).

[4]R. H. Dalitz and S. F. Tuan, Ann. Phys. (to be published).

[5]R. H. Dalitz and S. F. Tuan, Phys. Rev. Letters 2, 425 (1959).
[6]D. H. Stork and D. J. Prowse (private communication).
[7]Burrowes, Caldwell, Frisch, Hill, Ritson, and Schluter, Phys. Rev. Letters 2, 117 (1959).
[8]Nordin, Rosenfeld, Solmitz, Tripp, and Watson, Bull. Am. Phys. Soc. Ser. II, 4, 288 (1959); Eberhard, Rosenfeld, Solmitz, Tripp, and Watson, Phys. Rev. Letters 2, 312 (1959); A. H. Rosenfeld, Bull. Am. Phys. Soc. Ser. II, 3, 363 (1958).
[9]Cork, Lambertson, Piccioni, and Wenzel, Phys. Rev. 106, 167 (1957).
[10]Alvarez, Eberhard, Good, Graziano, Ticho, and Wojcicki (private communication of preliminary results).
[11]The determination made in reference 2 leads to $p_\Lambda X_\Lambda = 0.1 \pm 0.2$. This would suggest that both $p_\Lambda X_\Lambda$ and $p_\Sigma X_\Sigma$ are small, not only their "average."
[12]This quantity has also been discussed recently by P. T. Matthews and A. Salam, Phys. Rev. Letters 2, 226 (1959); J. D. Jackson and H. W. Wyld, Jr., Phys. Rev. Letters 2, 355 (1959).

MAGNITUDE OF RENORMALIZATION CONSTANTS*

S. G. Gasiorowicz
Lawrence Radiation Laboratory, University of California, Berkeley, California

and

D. R. Yennie and H. Suura
University of Minnesota, Minneapolis, Minnesota
(Received May 11, 1959)

Five years ago Källén[1] published a proof that at least one of the renormalization constants in quantum electrodynamics is infinite. A subsequent more refined analysis[2] led to the conclusion that it is the electron wave function renormalization which diverges. Recently Johnson[3] raised the question whether a formulation of the theory in another gauge might not alter this result. We have re-examined this problem and arrive at the conclusion that (a) Källén's proof is not conclusive, (b) the renormalization constants could be finite under rather special circumstances, and (c) the question of gauge invariance is quite irrelevant to this problem. We begin with a brief review of Källén's work on this problem.[4] It can be shown that the charge renormalization $(1-L)^{-1/2}$ can be expressed in the form

$$(1-L)^{-1} = 1 + \bar{\Pi}(0), \tag{1}$$

where by definition

$$\bar{\Pi}(Q^2) = P\int_0^\infty da\, \frac{\Pi(-a)}{a+Q^2}, \tag{2}$$

and

$$\Pi(Q^2) = -\frac{V}{3Q^2}\sum_{\substack{(n)\\ p^{(n)}=Q}} \langle 0|j_\mu|n\rangle\langle n|j_\mu|0\rangle; \tag{3}$$

P denotes the principal value. Källén shows that in spite of the indefinite metric associated with quantum electrodynamics, $\Pi(-a)$ is positive for positive a, and that therefore a lower limit to the integral in Eq. (1) can be obtained by considering any subset of eigenstates of the Hamiltonian. The simplest of these are the states consisting of an incoming electron-positron pair. Our task therefore is to study the high-energy limit of the matrix element $\langle 0|j_\mu|p,p'\rangle$.

By the use of reduction formulas[1,5] one may obtain the following compact expression for the matrix element in question:

$$\langle 0|j_\mu|p,p'\rangle = \langle 0|j_\mu^{(0)}|p,p'\rangle[1 - \bar{\Pi}(Q^2) + \bar{\Pi}(0) - i\pi\Pi(Q^2) + F_1(Q^2) - F_1(0)] + iQ_\nu\langle 0|m_{\mu\nu}^{(0)}|p,p'\rangle F_2(Q^2), \tag{4}$$

$$Q = p + p', \tag{4a}$$

where

$$N^2\bar{u}(p')\{\int\int d^4x d^4y\, e^{i(p'x+py)}[\theta(-x)\theta(-y)\langle 0|\{[j_\mu(0),\bar{f}(x)],f(y)\}|0\rangle - \theta(-y)\theta(y-x)\langle 0|\{[j_\mu(0),f(y)],\bar{f}(x)\}|0\rangle]\}v(p)$$

$$= F_1(Q^2)\langle 0|j_\mu^{(0)}|p,p'\rangle + iF_2(Q^2)Q_\nu\langle 0|m_{\mu\nu}^{(0)}|p,p'\rangle. \tag{5}$$

In terms of Källén's notation, $F_1(Q^2)$ and $F_2(Q^2)$ are given by[6]

$$F_1(Q^2)+F_2(Q^2)=\bar{R}^{\mathrm{reg}}(Q^2)+i\pi R^{\mathrm{reg}}(Q^2),$$
$$F_2(Q^2)=\bar{S}^{\mathrm{reg}}(Q^2)+i\pi S^{\mathrm{reg}}(Q^2). \quad (6)$$

Here N stands for the electron wave function renormalization constant and $f(x)$ is the fermion source density. The matrix elements with the superscript (0) are the lowest order perturbation expressions for the current and the spin. In this notation Ward's identity takes the form

$$F_1(0)=(1-N^2)/(1-L). \quad (7)$$

To estimate the behavior of this matrix element as $-Q^2\to\infty$, Källén argues as follows: let us assume that all the renormalization constants are finite. In particular this means that $\bar{\Pi}(0)$ is finite, so that $\Pi(-a)$ goes to zero fast enough as $a\to\infty$. For any reasonable[7] function $\Pi(-a)$, this implies that the same must hold for $\bar{\Pi}(Q^2)$. The problem therefore reduces to a study of the functions $F_1(Q^2)$ and $F_2(Q^2)$.

Källén originally examined the left-hand side of Eq. (5), and reduced it to an integral of the form

$$\int_{-\infty}^{\infty}\int dx\,dy\,\frac{F_\mu(\vec{p},x;\vec{p}',y)}{[x+y-(p_0+p_0'-i\epsilon)][y-p_0'+i\epsilon]}. \quad (8)$$

As $(p_0+p_0')\to\infty$ this integral need not vanish, because $\vec{p}$ and $\vec{p}'$ also vary when p_0 and p_0' vary (since $\vec{p}^2+m^2=p_0^2$; $\vec{p}'^2+m^2=p_0'^2$). The integral would vanish if the function $F_\mu(\vec{p},x;\vec{p}',y)$ approached zero as $\vec{p},\vec{p}'\to\infty$ but in general there is no way of establishing this. As the vanishing of this integral was an important part of Källén's original proof, we conclude that the proof is inconclusive.

A more constructive discussion can be given if we abandon a certain degree of generality, and assume that the functions $F_1(Q^2)$ and $F_2(Q^2)$ are analytic in the $-Q^2$ plane except for a cut along the positive real axis.[8] The behavior of the functions $F_1(Q^2)$, $F_2(Q^2)$ for large values of $-Q^2$ is unknown. The locality of the theory demands that they do not increase faster than a polynomial. In general we can therefore write

$$F_i(Q^2)=\frac{1}{\pi}\int_0^\infty da\,\frac{\rho_i(-a)}{a+Q^2-i\epsilon}+p_n^{(i)}(Q^2), \quad (9)$$

or, more relevantly,

$$\bar{R}^{\mathrm{reg}}(Q^2)=\bar{R}^{\mathrm{reg}}(0)+(+Q^2)\bar{R}^{\mathrm{reg}'}(0)+\ldots+\frac{(+Q^2)^{n-1}}{(n-1)!}\bar{R}^{\mathrm{reg}(n-1)}(0)$$
$$+(-Q^2)^n P\int_0^\infty \frac{da}{a^n}\,\frac{R^{\mathrm{reg}}(-a)}{a+Q^2}, \quad (10)$$

and $\bar{S}^{\mathrm{reg}}(Q^2)$ is given by a similar expression. The polynomials are seen to be real. It is important to note that the degree of the polynomial cannot be smaller but may be larger than the power necessary to make the integral over the imaginary part convergent.[9] Let us now discuss a few possibilities. We shall talk about R but the analysis applies equally well to S.

Case I. $n=0$.—In this case we have

$$\bar{R}^{\mathrm{reg}}(Q^2)=P\int_0^\infty da\,\frac{R^{\mathrm{reg}}(-a)}{a+Q^2}. \quad (11)$$

This implies that

$$\bar{R}^{\mathrm{reg}}(Q^2)-\bar{R}^{\mathrm{reg}}(0)=-\int_0^\infty \frac{da}{a}\,\frac{R^{\mathrm{reg}}(-a)}{1+a/Q^2}, \quad (12)$$

so that

$$\lim_{-Q^2\to\infty}\bar{R}^{\mathrm{reg}}(Q^2)=\bar{R}^{\mathrm{reg}}(0)-\int_0^\infty\frac{da}{a}R^{\mathrm{reg}}(-a)=0. \quad (13)$$

If $\bar{S}^{\mathrm{reg}}(Q^2)$ has the same behavior, we obtain [using Eq. (7)] the result

$$\langle 0|j_\mu|p,p'\rangle\to\frac{N^2}{1-L}\langle 0|j_\mu^{(0)}|p,p'\rangle, \quad (14)$$

i.e., the matrix element approaches the unrenormalized Born approximation, and a contradiction with the initial assertion that all renormalization constants are finite is obtained because $\Pi(-a)\to\mathrm{const}\ N^2/(1-L)^2$ as $a\to\infty$. This corresponds to the case discussed by Källén.

Case II. $n=1$.—In this case we have

$$\bar{R}^{\mathrm{reg}}(Q^2)=\bar{R}^{\mathrm{reg}}(0)+(-Q^2)\int_0^\infty \frac{da}{a}\,\frac{R^{\mathrm{reg}}(-a)}{a+Q^2}. \quad (15)$$

If the integral

$$\int_0^\infty \frac{da}{a}\,R^{\mathrm{reg}}(-a) \quad (16)$$

converges, we have

$$\lim_{-Q^2\to\infty}\bar{R}^{\mathrm{reg}}(Q^2)=\bar{R}^{\mathrm{reg}}(0)-\int_0^\infty da\,\frac{R^{\mathrm{reg}}(-a)}{a}. \quad (17)$$

The difference between this case and the one discussed before is that the right side of Eq. (17) need not vanish. In particular, if this constant is such as to make $\langle 0|j_\mu|p,p'\rangle \to 0$ as $-Q^2\to\infty$, (this necessitates that there be no polynomial in the S^{reg} equation), then no conclusion as to the magnitude of the renormalization constants can be drawn from the electron-positron pair states alone. This is the only possible exception to Källén's result. If the integral in Eq. (16) does not converge, $\bar{R}^{\mathrm{reg}}(Q^2)$ will diverge in the limit $-Q^2\to\infty$ for any reasonable function $R^{\mathrm{reg}}(-a)$. Thus except for the special case noted above, the integral in Eq. (2) will always diverge (this is even more easily seen if $n=2, 3, \ldots$).[10] The special case corresponds to the "no-subtraction" spectral representations for the matrix elements of the current operator discussed recently.[11]

It is clear from this discussion that developing the theory in another gauge cannot make any difference. We have used gauge invariance only in the particular form in Eq. (5), but the scalar functions $F_1(Q^2)$ and $F_2(Q^2)$ can be discussed independently of whether the photon has a mass or not. This fact will affect the low-frequency behavior of the functions [e.g., Eq. (7)] but cannot make any difference in the high-energy region. Our negative conclusion applies by the same token to all renormalizable theories: one cannot prove that the renormalization constants are infinite, because there is always one possibility that a special cancellation occurs.

We may ask whether it is possible that the renormalization constants are indeed finite. For this to be true all matrix elements in Eq. (3) must have the "no-subtraction" form. Actually this is not enough, because even though every contribution to the integral $\bar{\Pi}(0)$ will be finite, it is not necessarily true that the sum of these contributions converges. In fact, since $\langle 0|j_\mu|n\rangle$ could not be expected to approach zero until $a>>(E_0N)^2$ (where E_0 is some fixed energy and N is the number of particles in the state), a very nonuniform series is involved, and counterexamples to convergence can readily be constructed.[12] Thus unless field theory has some regularity properties quite unsuspected from perturbation theory, the renormalization constants will be infinite.

In our discussion we have given up a certain amount of generality by working with the unproved spectral representations. Should these turn out to be wrong, there may be two possibilities: If the violation of analyticity is limited to a finite part of the plane (e.g., longer cuts), then the domination of the high-energy behavior by a polynomial is unaltered and our discussion in terms of spectral representations needs only minor modifications. This would not be true, however, if regions of nonanalyticity extend to infinity (e.g., cuts not indicated by perturbation theory such as those which might arise from complex "ghosts").

This work was started while one of us (S.G.) was visiting the Department of Physics at the University of Minnesota. He is grateful to Professor A. O. Nier and the Department of Physics for their kind hospitality.

---

*This research was supported in part by the U. S. Atomic Energy Commission.

[1]G. Källén, Kgl. Danske Videnskab. Selskab, Mat.-fys. Medd. 27, No. 12 (1954).

[2]G. Källén, CERN Report T/GK/3 (unpublished).

[3]K. A. Johnson, Phys. Rev. 112, 1367 (1958).

[4]G. Källén, Handbuch der Physik (Springer-Verlag, Berlin, 1958), Vol.V, Part 1, pp. 358-363.

[5]Lehmann, Symanzik, and Zimmermann, Nuovo cimento 2, 425 (1955).

[6]The superscript "reg" refers to a particular prescription for calculating these functions. This complication introduced by the necessity of adiabatic decoupling is of no import in the following argument, so that we will not take the space to discuss it, but rather refer the interested reader to reference 4.

[7]This point, summarized here so cavalierly, is discussed in detail in the appendix of reference 1.

[8]This property has been established to all orders in perturbation theory [Y. Nambu, Nuovo cimento 6, 1064 (1957)], and under certain circumstances, unrealistic for this theory, independent of perturbation theory [see Bremermann, Oehme, and Taylor, Phys. Rev. 109, 2178 (1958)].

[9]This was overlooked in an analysis of the problem

in terms of spectral representations given by G. Källén in Proceedings of the CERN Symposium on High-Energy Accelerators and Pion Physics, Geneva, 1956 (European Organization of Nuclear Research, Geneva, 1956), Vol. II, p. 187.

[10]Källén remarked in his paper that a divergence of $\lim_{\zeta^2\to\infty}\bar{R}^{\mathrm{reg}}(\zeta^2)$ as well as its vanishing will lead to infinite renormalization constants, but he overlooked the only possible counterexample.

[11]S. Drell and F. Zachariasen, Phys. Rev. 111, 1727 (1958); M. Gell-Mann and J. Mathews (to be published).

[12]See reference 9. The infrared divergence in quantum electrodynamics implies a nonuniform behavior of this type for the series in Eq. (3). According to arguments of Bloch-Nordsieck type, the matrix element for producing an electron pair plus any finite number of photons should vanish. Nevertheless, the sum over all numbers of photons should yield a finite result.

ERRATUM

PRECISE DETERMINATION OF THE MUON MAGNETIC MOMENT. R. L. Garwin, D. P. Hutchinson, S. Penman, and G. Shapiro [Phys. Rev. Letters 2, 213 (1959)].

It has recently been[1] established that the theoretical uncertainties in the lower limit of the muon mass from mesonic α-rays[2] are much smaller than previously supposed.[3] We therefore undertook a more careful search into possible systematic errors in our moment measurement. We have found and corrected one serious source of error. A method of testing was then designed which simulated the experimental situation and permitted the establishment of an upper limit on the remaining systematic errors in the electronic equipment of 1×10^{-5} or almost an order of magnitude less than our stated accuracy.

The circuit design used in this experiment has been intended to be completely aperiodic, i.e., should introduce no systematic errors in a frequency measurement (as opposed to a random error in measuring each particle due to the finite time resolution). Preceding the analysis equipment, however, is the "zero-crossing detector" which does the fast timing of both the muon and decay electron pulses. To circumvent the difficulties of long-term transit-time variations in phototubes, a single counter and timing circuit was used for both muon and electron pulses. This raises the possibility that the recovery of the phototube or fast timing circuit after the passage of a muon pulse will alter the apparent time of the electron. This effect is important only if it varies during the measurement interval. In particular, if electrons immediately following the muon pulse are delayed more than those coming later in the gate interval, the net effect is to make the apparent muon precession frequency appear larger than it actually is. Unfortunately, we observed on our circuit diagram a one-microsecond time-constant at the grid following the distributed zero-crossing amplifier. Had this time-constant been 50 millimicroseconds or >20 microseconds, no error would have been incurred, but we estimated the maximum error caused by it as $\sim4\times10^{-4}$.

Since there is available to us no measuring equipment that approaches the experimental apparatus in resolving ability, the existence of the postulated effect was established by inserting two pulses with fixed cable delay of 2 microseconds into the equipment. The first was analyzed by the circuitry as a muon pulse and the other as an electron pulse. The equipment output was displayed on a pulse-height analyzer as previously described. A third pulse which had been traveling in a long cable was then inserted between the "muon" and "electron" pulses. Its retarding effect on the electron timing could then be seen clearly. The fast timing circuit was then altered (by changing the offending time-constant) so that the effect was no longer apparent.

After modifying the circuitry, the following rigorous test was made to establish that there remained no other such sources of error. We built a source of random pulses whose probability of occurrence above a fixed threshold oscillated in time at a known rate. This consisted of a 6810A photomultiplier viewing a plastic scintillator which was irradiated with β rays. The focusing voltage of the tube was modulated at or near the frequency of the reference oscillator of the measuring apparatus, (86.2 Mc/sec). A 30-volt peak-to-peak signal was sufficient to achieve nearly 100% modulation of the output pulses. The slow coincidence circuits select the first pulse above a threshold after a dead time as a μ pulse and the next pulse as an electron. Since the probability of occurrence of these two

pulses is strongly correlated at the reference oscillator frequency, the pulse-height analyzer display of the electron phase modulo the reference oscillator period will show a structure very similar to that for a precessing muon [although not restricted to the form $(1+a\cos\phi)$]. The actual experiment consists of comparing the phase distribution of electrons that come soon after a muon stop with those that come later. In the test, two groups of modulated random pulses were displayed, those that corresponded to "early" electrons and those that corresponded to "late" ones.

Figure 1 shows the actual pulse-height analyzer display. On the left is the distribution in phase of the early "electrons" and on the right that of the late "electrons." Ideally, the distributions will be identical in shape. A treatment of the test data by the same Fourier analysis used in the experiment shows no phase shift corresponding to a frequency error larger than 10^{-5}.

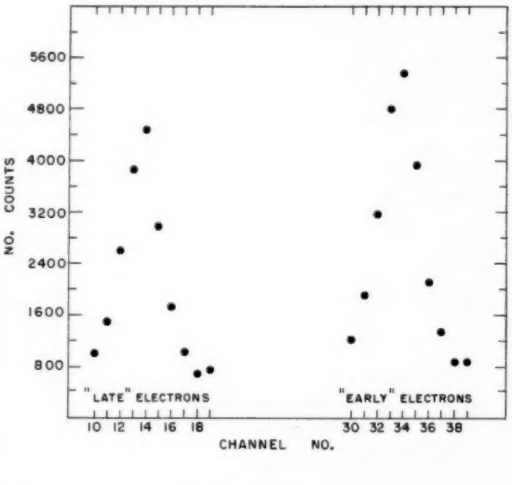

FIG. 1. Pulse-height analyzer display of the modulated counter test.

The measurement of the muon moment was then repeated using a smaller bromoform target, with the equipment modified and tested as above.

Expressing the final result in terms of the ratio of the resonance frequency of the muon to that of a proton in the same magnetic field, we obtained $f_\mu/f_p = 3.1834 \pm 0.0002$. Combining this value for the moment with the value for the lower limit on the mass, as discussed in reference 1, we obtain for the g factor $g_\mu > 2(1.00122 \pm 0.00008)$. This limit is seen to be in agreement with the theoretical prediction $g = 2(1.00116)$. The experiment to determine the absorption coefficient of lead for the $3D$ - $2P$ transition in μ-mesonic phosphorus is being repeated. Parratt[4] has pointed out that the width of the lead K edge should enable one to obtain either a more accurate mass value or a considerably different lower limit.

---

[1]A. Petermann and Y. Yamaguchi, Phys. Rev. Letters 2, 359 (1959).

[2]Koslov, Fitch, and Rainwater, Phys. Rev. 95, 291 (1954).

[3]Cohen, Crowe, and DuMond, Fundamental Constants of Physics (Interscience Publishers, New York, 1957).

[4]L. G. Parratt (private communication).

ABSTRACTS

In this section are printed the abstracts of Articles that have been forwarded to The American Institute of Physics for publication in THE PHYSICAL REVIEW. In quoting information obtained from this section before the appearance of the corresponding Article, reference should be made to "Physical Review (to be published)" rather than to this Journal.

MEASUREMENT OF DISTANCE IN GENERAL RELATIVITY. E. Newman, University of Pittsburgh, Pittsburgh, Pennsylvania, and J. N. Goldberg, Aeronautical Research Laboratory, Wright-Patterson Air Force Base, Ohio (Received January 28, 1959).

An intrinsic geometrical definition of distance is presented in terms of the geodesic deviation of null rays. This definition is applied to the Schwarzschild solution and the homogeneous and isotropic cosmological solutions of the Einstein field equations. In the former case the distance is proportional to the radial coordinate of the standard metric, and in the latter case it yields the usual cosmological distance. However, since the geodesic deviation is a scalar, the results are independent of the particular metric used. The relationship of this definition to observation is discussed and it is concluded that it agrees with the astronomical definition of luminosity distance.

POSSIBLE SUPERFLUIDITY OF A SYSTEM OF STRONGLY INTERACTING FERMIONS. L. N. Cooper,* R. L. Mills, and A. M. Sessler, Ohio State University, Columbus, Ohio (Received January 30, 1959).

The possible superfluidity of a system of strongly interacting fermions is investigated on the assumption that an adequate description of the system in its "normal" state is given by independent fermions in a momentum-dependent potential. On the basis of this assumption we have investigated whether a correlated wave function of the form used by Bardeen, Cooper, and Schrieffer, minimizes the ground-state energy. The nonzero terms in the expectation value of the Hamiltonian contain the modified kinetic energy and the full two-body potential between the fermion pairs. An integral equation is obtained in configuration space for the correlation function between pairs. This integral equation is meaningful even for potentials with hard cores, and a nonzero solution implies the existence of a superfluid state. A variational method is devised which provides a criterion for superfluidity and a lower bound for the transition temperature into the superfluid state. We find that a repulsive hard core does not in principle forbid the existence of a superfluid state, but whereas in the absence of a hard core an attractive two-body potential always leads to a superfluid state at sufficiently low temperatures, in the presence of a repulsive core there appears to be a critical strength of attraction needed to form a superfluid state. When the variational principle is applied to liquid He^3 or to nuclear matter, it is found for a wide class of trial functions that the system does not become a superfluid.

*Present address: Brown University, Providence, Rhode Island.

STATIONARY NONEQUILIBRIUM GIBBSIAN ENSEMBLES. J. L. Lebowitz, Department of Physics, Stevens Institute of Technology, Hoboken, New Jersey (Received January 6, 1959).

The general theory of a Gibbs ensemble representing a system in contact with its surroundings is applied to several concrete situations of interest. By an appropriate choice of heat reservoirs a simply modified Liouville equation is found to describe a heat conducting system. The stationary nonequilibrium Γ-space ensembles which describe such a system are found explicitly for some cases. In the simplest cases these ensembles turn out to be canonical with a temperature that is a weighted average of the reservoir temperatures. For other systems, such as Brownian particles inside a fluid whose temperature is not uniform, we find the stationary ensemble to terms linear in the temperature gradient. From this we are led to discuss ensembles that will approximately represent an arbitrary heat conducting fluid. A more general proof than previously given for the asymptotic approach of the Γ-space distribution to its stationary value is also presented.

VELOCITIES OF SPUTTERED ATOMS. G. K. Wehner, Mechanical Division of General Mills, Incorporated, Minneapolis, Minnesota (Received January 5, 1959).

The average velocity of atoms sputtered from metal surfaces by normally incident Hg^+ ions in

a low-pressure discharge was found to be unexpectedly high, corresponding to kinetic energies of the order of 10 to 30 ev.

By sputtering atoms from a (110) surface of a face-centered cubic crystal, rather dense atomic beams ejected normal to the target surface can be obtained. The fact that the kinetic energies of the atoms are more than 100 times higher than thermal evaporation energies makes such a beam interesting.

OPTICAL PROPERTIES OF NICKEL OXIDE. R. Newman and R. M. Chrenko, General Electric Research Laboratory, Schenectady, New York (Received January 28, 1959).

The absorption spectrum of single-crystal NiO has been measured in the 0.1 to 6 ev range and its reflectivity spectrum from 0.025 to 10 ev. The absorption spectrum shows a series of lines from 1 to 3.5 ev arising from internal transitions of the Ni ion. A continuous background absorption occurs in the range from 0.1 to 3.5 ev, whose magnitude increases with impurity concentration. The absorption coefficient rises steeply above 3.5 ev and reaches a value of 10^6 cm^{-1} at and above 4 ev. An absorption line at 0.24 ev is found to be temperature sensitive in both intensity and frequency in the range above 300°K. Its behavior suggests that it is connected with the antiferromagnetic ordering. The reststrahlen spectrum was observed with the following parameters: high- and low-frequency dielectric constants 5.4 and 12, respectively; energies of longitudinal and transverse optical mode vibrations are 0.076 and 0.044 ev, respectively.

STRESS AND TEMPERATURE DEPENDENCE OF THE PARAMAGNETIC RESONANCE SPECTRUM OF NICKEL FLUOSILICATE. W. M. Walsh, Jr., Gordon McKay Laboratory, Harvard University, Cambridge, Massachusetts (Received January 19, 1959).

The crystalline field splitting of the ground-state spin triplet of divalent nickel in the fluosilicate, $NiSiF_6 \cdot 6H_2O$, has been measured at room temperature as a function of hydrostatic pressure to 10 000 kg/cm² and uniaxial stress to 150 kg/cm². The anisotropic compressibility and thermal expansion of this trigonal crystal have also been determined. When these data are combined with the known variation of the splitting with temperature, the dependence of the splitting on isothermal unit-cell geometry and on temperature at constant unit-cell dimensions is calculated. The splitting is found to be independent of volume within experimental error but proves to be quite sensitive to unit-cell shape. The deduced explicit temperature dependence is three times larger than that measured at atmospheric pressure. The magnitude and geometrical variation of the crystalline field splitting may be qualitatively understood using a static, ionic model of the $(Ni \cdot 6H_2O)^{2+}$ octahedral complex. A rather general analysis of the explicit temperature dependence indicates, however, that low-frequency lattice vibrations play a dominant role in determining the observed value of the splitting.

The resonance line widths are observed to increase monotonically and quite nonlinearly with increasing pressure. This broadening is discussed in terms of isotropic and anisotropic exchange interactions. In agreement with earlier conclusions of Ollom and Van Vleck, it is inferred that the two mechanisms are of comparable importance in this paramagnetic salt.

PRESSURE DEPENDENCE OF THE PARAMAGNETIC RESONANCE SPECTRA OF TWO DILUTE CHROMIUM SALTS. Walter M. Walsh, Jr., Gordon McKay Laboratory, Harvard University, Cambridge, Massachusetts (Received January 19, 1959).

The pressure dependences of the spin-Hamiltonian parameters of trivalent chromium in ammonium aluminum alum and potassium cobalticyanide have been measured up to 10 000 kg/cm² near room temperature. In the case of the alum, runs have been made at 0, 24, and 50°C. The g-value remains unchanged while the crystalline field splitting, δ, increases by ~30%, $(\partial\delta/\partial P)_T$ decreasing with rising temperature and pressure. The compressibility of this cubic crystal is known but it was necessary to measure the thermal expansion coefficient in order to complete the equation of state. The crystal contracts when heated: $(\partial \ln V/\partial T) = -8.1 \times 10^{-5}(°C)^{-1}$ in the interval (0, 50°C). This anomalous result is consistent with the temperature dependence of the compressibility found by Bridgman and may correlate with the known low-temperature phase transition of the alum. Using the empirically determined equation of state, the crystalline field splitting variations are converted to isothermal volume and explicit tem-

perature dependences. These are discussed in terms of static and dynamic crystalline fields, but are not satisfactorily explained. This failure is attributed to inhomogeneous internal deformation of the unit cell as a function of stress and temperature.

A room temperature run on the covalent chromicyanide shows the g-value as well as the principal splitting parameter, D, to be nonmonotonic functions of pressure. The rhombic splitting parameter, E, increases quadratically with pressure. Since no attempt was made to determine the crystalline equation of state, the resonance data are only qualitatively discussed. It is difficult to reconcile the results with the equivalent crystalline-field model of the chromicyanide complex.

THERMOLUMINESCENCE AND F CENTERS. I. THEORY. G. Bonfiglioli, P. Brovetto, and C. Cortese, Istituto Elettrotecnico Nazionale, Turin, Italy (Received December 22, 1959).

After a critical review of the present state of knowledge about thermoluminescence, which is shown to be rather unsatisfactory, the present paper considers in detail thermoluminescence arising from thermal bleaching of F centers in irradiated alkali halides. A model is proposed, which leads to second-order kinetics for the radiative recombination of F centers and Luminophor centers (L_i) which must be present in the crystal lattice.

The formalism concerned is developed, showing how activation energy and probability factors can be determined.

A quantitative argument is presented, concerned with the nature of L_i centers and suitable for the identification of their ground levels within the forbidden crystal band.

THERMOLUMINESCENCE AND F CENTERS. II. EXPERIMENTAL. G. Bonfiglioli, P. Brovetto, and C. Cortese, Istituto Elettrotecnico Nazionale, Turin, Italy (Received December 22, 1959).

The theoretical treatment developed in the previous paper is applied to interpret experiments of thermoluminescence associated with the presence of F centers in irradiated alkali halides. Some important features of the technique employed are discussed—in particular, the procedure used to establish the proper wavelength belonging to the light emitted in each glow peak.

The energy of activation and the probability factors of the phenomenon are determined, and it is shown that the second and third peaks of thermoluminescence are probably due to direct radiative recombination of F centers with V_2 and V_3 centers. On the other hand, it is still uncertain whether the first peak is associated with V_1 centers or with Cu ion impurities in the NaCl lattice.

TEMPERATURE DEPENDENCE OF THE ELECTRICAL PROPERTIES OF BISMUTH-ANTIMONY ALLOYS. A. L. Jain, Department of Physics and Institute for the Study of Metals, University of Chicago, Chicago, Illinois (Received January 22, 1959).

The electrical resistivity ρ and Hall effect R of zone-levelled single crystals of Bi-Sb alloys have been measured in the temperature range from 4.2°K to 300°K. $\mathrm{Log}(\rho/\rho_{300})$ vs $1/T$ curves suggest thermal activation of carriers in the concentration range from 5% to 40% Sb in the temperature range from 25°K to 100°K; approximate activation energies have been inferred from their slopes. The activation energy appears to have a maximum at a concentration near 12%. Some anomalies have been observed in the behavior of ρ and R on both sides of this concentration at low temperatures. Lattice parameters for these alloys have also been measured for the entire range of solid solubility. Both a maximum and minimum in the c-axis lattice parameter vs concentration occur near the concentration at which anomalies appear in transport properties. These phenomena are discussed in terms of a simple band model proposed by Blount and Cohen.

ELECTRON BOMBARDMENT OF SILICON. D. E. Hill,* Purdue University, West Lafayette, Indiana (Received January 22, 1959).

The electrical properties of single-crystal silicon have been investigated as a function of bombardment with high-energy electrons. Measurements of the initial carrier removal rate were carried out at 273°K for samples with widely different carrier concentrations. A second type of experiment involved the measurement of the temperature dependence of resistivity and Hall

coefficient at different stages of the bombardment. Bombardment-produced energy levels were found at 0.03 ev, 0.17 ev, and 0.4 ev below the edge of the conduction band. With 4.5-Mev electrons, the introduction rates for these levels were 11 cm^{-1}, 0.5 cm^{-1}, and 0.05 cm^{-1}, respectively. Bombardment levels were also found at 0.05 ev and 0.3 ev above the edge of the valence band. With 4.5-Mev electrons, the introduction rates for these levels were 13 cm^{-1} and roughly 0.3 cm^{-1}, respectively. Irradiations with 700-kev electrons indicate the presence of the same defect levels with correspondingly smaller introduction rates. An investigation of the irradiation-induced mobility changes for degenerate p-type silicon indicates that, for each hole removed from the valence band, there appears a singly-charged scattering center. For near-degenerate n-type samples, the mobility changes are much smaller. Results of bombardments at 78°K and 10°K show that the net damage rate decreases by a factor of 2 or more when bombardments are carried out at lower temperatures.

*Now at Monsanto Chemical Company, Dayton, Ohio.

CYCLOTRON ABSORPTION IN METALLIC BISMUTH AND ITS ALLOYS. J. K. Galt, W. A. Yager, F. R. Merritt, and B. B. Cetlin, Bell Telephone Laboratories, Murray Hill, New Jersey, and A. D. Brailsford, Ford Motor Company, Dearborn, Michigan (Received January 21, 1959).

Extensive studies of cyclotron absorption in bismuth and in alloys of bismuth with tin and tellurium are reported. Experimentally, a microwave cavity arrangement is used which causes circularly polarized radiation to be incident on the sample. The samples are disks which form part of the cavity wall during the experiment. A signal is observed which is under most conditions proportional to the change in the fractional power absorption of the sample as the magnetic field applied to the sample is changed. The experiment is done with this field in the plane of the sample as well as normal to it. When the magnetic field is normal to the sample surface, circularly polarized radiation distinguishes between the effects of holes and of electrons. Making this distinction is complicated in bismuth by the anisotropy of the band structure, but it is not impossible. In pure bismuth the number of holes equals the number of electrons; alloying with tin increases the relative number of holes and alloying with tellurium increases the relative number of electrons. The experiments have been done both at 24 000 Mc/sec and at 72 000 Mc/sec. The tilted ellipsoid model of the electron band,[39] and the ellipsoid of revolution model for the hole band[33] have been used in a classical magneto-ionic theory of the effect which fits the data taken with magnetic field normal to the sample surface quite satisfactorily. The fit between this theory and the experiment is substantial confirmation of the validity of this band model, and makes it possible with the aid of certain infrared observations to determine the parameters of the model numerically. The mass parameters obtained for an ellipsoid in the electron band are $m_1 = 0.0088$, $m_2 = 1.80$, $m_3 = 0.02323$, $m_4 = \pm 0.16$, while those for the hole band are $M_1 = M_2 = 0.068$ and $M_3 = 0.92$. Comparison is made with similar results obtained from magneto-resistance and de Haas - van Alphen experiments. The experiments done with the magnetic field in the plane of the sample have been used to improve the accuracy of the determination of cyclotron resonance fields and therefore of the masses of the charge carriers, although no quantitative theory of the over-all shape of the signal under these conditions is available.

NUCLEAR MAGNETIC RESONANCE ACOUSTIC ABSORPTION IN KI AND KBr. D. I. Bolef and M. Menes, Physics Department, Westinghouse Research Laboratories, Beulah Road, Churchill Borough, Pittsburgh, Pennsylvania (Received January 28, 1959).

A technique for measuring the coupling between nuclear spins and the crystal lattice is described. The technique is one in which the absorption of energy by the nuclear spins from an acoustic wave is observed as an additional attenuation of the wave in the sample. This attenuation results in a change in the electric impedance of the piezoelectric transducer which drives the sample. The transducer impedance is made to control a Pound-Watkins type of spectrometer through an inverting and matching network.

Nuclear magnetic resonance acoustic absorption measurements were made on single crystal samples of KI and KBr using compressional waves in the [100] direction. These measurements yield a value of the electric quadrupole coupling $e^2 q_1 Q$, where q_1 is the first derivative with respect to strain of the electric field gradi-

ent at the nuclear position. For I^{127} in KI, $e^2q_1Q/h = 300$ Mc/sec. For Br^{81} in KBr, $e^2q_1Q/h = 82$ Mc/sec. Interpreting these results on the basis of a point charge model, we obtain values of the amplification factor γ. For KI^{127}, $\gamma = 38$. For KBr^{81}, $\gamma = 26$.

MILLIMETER-WAVE ROTATIONAL SPECTRUM OF NO IN THE $^2\pi_{3/2}$ STATE. Paolo G. Favero,* Anna M. Mirri, and Walter Gordy, Department of Physics, Duke University, Durham, North Carolina (Received January 26, 1959).

The $J = 3/2 - 5/2$ rotational transition of NO in the $^2\pi_{3/2}$ electronic state has been measured in the 1.17-mm wave region. Theory applied to these measurements combined with previous measurements on transitions of the ground $^2\pi_{1/2}$ state lead to the following values for the characteristic constants of $N^{14}O^{16}$: The rotational constant, $B_0 = 50\,848.42$ Mc/sec, and the spin-orbit coupling constant $A = 122.094$ cm^{-1}. The N^{14} nuclear magnetic couplings $a = 83.82$ Mc/sec, $b = 68.49$ Mc/sec, $c = -86.34$ Mc/sec, and the N^{14} nuclear quadrupole couplings $eQq_1 = -2$ Mc/sec, $eQq_2 = 22$ Mc/sec. The Λ-doubling constants $p_\Lambda = 175.15$ Mc/sec, $q_\Lambda = 1.15$ Mc/sec.

*Permanent address: Institute of Physical Chemistry, University of Padua, Padua, Italy.

SCATTERING OF ELECTRONS WITH AN INTRINSIC ELECTRIC DIPOLE MOMENT. B. Margolis, S. Rosendorff, and A. Sirlin, Department of Physics, Columbia University, New York, New York (Received January 19, 1959).

The elastic scattering of electrons with a small intrinsic electric dipole moment by spinless targets is discussed in the first and second Born approximations. For large momentum transfer ($q \sim 500$ Mev/c) the electric dipole moment contribution to the elastic scattering cross section becomes important even for values $\lambda \sim 10^{-3}$ (where λ is the electric dipole moment in units of e times the Compton wavelength of the electron). Thus, it seems likely that experiments involving high-energy electrons may give valuable information on the hypothetical existence and upper limit of the electric dipole moment interaction. It is also found, in the first Born approximation, that at sufficiently high energies the existence of the electric dipole moment can change the polarization of the electron beam from longitudinal to transverse. Finally, asymmetry effects which appear in the second Born approximation are discussed in detail. Detection of these asymmetries or of the spin correlation effects of the first Born approximation mentioned above would imply the noninvariance of the electromagnetic interactions under parity and time reversal.

IONIZATION PRODUCED BY ATOMIC COLLISIONS AT kev ENERGIES. II. A. Russek and M. Tom Thomas,† The University of Connecticut, Storrs, Connecticut (Received January 26, 1959).

The electron evaporation model, presented in a previous paper, of the collision-ionization process that occurs when atoms collide at high energies is extended and improved. In addition, an alternate model of the process is considered. This second model treats the collision-ionization process as being due to direct knock-outs of the electrons by violent electron-electron collisions. It is found that this model cannot be made to account for the data, thereby lending additional support to the assumptions inherent in the evaporation model.

†Present address: Physics Department, Brown University, Providence, Rhode Island.

DECAY OF Zn^{61}. James B. Cumming, Chemistry Department, Brookhaven National Laboratory, Upton, New York (Received January 26, 1959).

The 88.5-sec activity observed among the products of the bombardment of nickel with 16-Mev alpha particles is assigned to Zn^{61}. The decay properties of this isotope have been investigated by beta and gamma scintillation techniques and by coincidence measurements. Approximately 80% of the Zn^{61} beta decays are to the ground state of Cu^{61}. The positron end point is 4.38 Mev which corresponds to a logft of 5.2 and to an allowed transition. Gamma rays of 0.48, 0.98, and 1.64 Mev have been observed with intensities of 11.2, 2.9, and 6.2 percent of the positrons. There is some indication of gamma rays at 0.69 Mev and possibly other energies in lower intensities. A decay scheme is proposed which incorporates the above data into an existing level scheme of Cu^{61} derived from (p,γ) experiments.

DECAY OF Co^{56} and Mn^{56}. P. Kienle,* Brookhaven National Laboratory, Upton, New York, and R. E. Segel, Aeronautical Research Laboratory, Dayton, Ohio (Received January 23, 1959).

Scintillation spectroscopy and coincidence techniques have been used to measure the decay schemes of Mn^{56} and Co^{56}, both of which decay to Fe^{56}. Gamma-ray energies and intensities were measured with the aid of a large NaI(Tl) crystal and a three-crystal pair spectrometer. Combining these data with γ-γ and β-γ coincidence measurements, decay schemes are constructed for both nuclei. The logft values for the various β branches are examined and all of the β branches from Co^{56} show higher than usual logft values for allowed transitions. The Co^{56} decay to the second excited state in Fe^{56} is inhibited by a factor of ~10^3 while the Mn^{56} decay to the same state is inhibited even further. The allowed Mn^{56} decay to the first excited state in Fe^{56} is also slowed, while Mn^{56} decays to higher states proceed at more normal speeds. A discussion of the β-transition rates in terms of shell model assignments is given, from which some explanation of the β rates results. The high inhibitions of the decays to the second excited state in Fe^{56} do not seem susceptible to a simple shell model explanation.

*Permanent address: Technische Hochschule, München, Germany.

THEORY OF DIRECT-INTERACTION INELASTIC SCATTERING. Norman K. Glendenning, Lawrence Radiation Laboratory, University of California, Berkeley, California (Received January 12, 1959).

A model is proposed for the description of direct-interaction inelastic scattering in which resolved final states of the target nucleus are observed. This model is expected to be useful in deducing the spins and parities of the excited states of nuclei.

INTERNAL CONVERSION IN Li^7. R. D. Leamer,* Carnegie Institute of Technology, Pittsburgh, Pennsylvania (Received January 26, 1959).

Upper and lower limits have been placed on the internal conversion coefficient of the 479-kev excited state of Li^7. The limits are 16.6×10^{-7} and 5.8×10^{-7}, yielding a multipole mixture of (100% - 55%) $M1$ and (0% - 45%) $E2$ for the gamma ray. Photoelectron emission from the source and the background counting rate impose the principal limitations on the accuracy of the experiment.

*Now at Westinghouse Reactor Evaluation Center, Ruffsdale, Pennsylvania.

FINITE-RANGE DELTA-FUNCTION POTENTIAL. M. Bolsterli, School of Physics, University of Minnesota, Minneapolis, Minnesota (Received January 26, 1959).

The long-range part of the nucleon-nucleon potential is approximated by the finite-range delta-function potential:

$$V_{\text{singlet even}} = V_1\delta(r-a),$$

$$V_{\text{triplet even}} = V_3\delta(r-a) + V_T\delta(r-a)S_{12}.$$

The best fit to the low-energy n-p data is obtained when V_3 is repulsive and V_T is large. The uncertainties in the V's are of the order of 10%.

NEUTRON ACTIVATION AT 195 kev. W. S. Lyon and R. L. Macklin, Oak Ridge National Laboratory, Oak Ridge, Tennessee (Received January 19, 1959).

Neutron activation cross sections have been measured at 195 kev (50 kev energy spread) for about 30 isotopes, by absolute gamma counting. Absolute calibration was effected by determining the RdTh-D_2O neutron source strength with a flat-response 4π graphite-sphere neutron counter, and irradiating indium in the form of a thin, uniform, spherical shell. The results are in general agreement with those of other investigators. The 25-kev activation cross sections of Pr^{141}, Zr^{94}, Ga^{71}, and Pd^{108} have been revised.

EXPERIMENTAL DETERMINATION OF THE NONMAGNETIC NEUTRON-ELECTRON INTERACTION. E. Melkonian, B.M. Rustad, and W. W. Havens, Jr., Columbia University, New York, New York, and Brookhaven National Laboratory, Upton, New York (Received January 21, 1959).

A precise set of measurements of the variation of the transmission of liquid bismuth with energy in the range 0.1 ev to 10 ev has been made. After correction for neutron capture, Doppler effect, and liquid diffraction, a value of $V_0 = -4.34 \pm 0.14$

kev is obtained as the depth of the potential well having a radius equal to the classical electron radius. This gives -0.26 ± 0.14 kev as the strength of the intrinsic neutron-electron interaction after subtraction of the Foldy term.

SELF-INDICATION RESONANCE NEUTRON EXPERIMENTS WITH SILVER (2-200 ev). James E. Draper,* Brookhaven National Laboratory, Upton, New York, and Yale University, New Haven, Connecticut (Received June 5, 1958; revised manuscript received January 19, 1959).

A method is given for the area analysis of resonance neutron self-indication experiments including consideration of the effects of energy resolution, wing area, energy dependence of neutron flux, v^{-1} term in the capture cross section, potential scattering, detailed treatment of Doppler broadening, and scattering events preceding capture in the detector sample. The essential result of this measurement alone is nearly the peak total cross section of the Doppler-broadened resonance (viz., $\sigma_{t0}\psi_0/\psi_0^{1/3}$) which, with approximate knowledge of Γ, gives σ_{t0}. Experimental results are given for eleven resonances below 200 ev in silver, and they are compared with results obtained by other methods.

*Present address: Yale University, New Haven, Connecticut.

METHOD FOR THE DETERMINATION OF STRANGE PARTICLE PARITIES AND COUPLING CONSTANTS. John G. Taylor, Christ's College, Cambridge, England (Received January 7, 1959).

A general method for determining parities and coupling constants of strange particles directly from angular distributions is described. A reasonably detailed discussion is given of what can be gleaned by this method from the processes of associated production by photons on nucleons and by pions on protons, and from the absorption and charge exchange scattering of charged K-mesons on nucleons. If the relative parity of K^+ to K^0 is odd then strange particle parities may be obtained from all these processes, while if this relative parity is even only photoproduction processes may be used to determine parities. Once these relative parities have been determined then coupling constants may be obtained from all the processes.

HYPERFINE SPLITTING EFFECTS IN THE CAPTURE OF POLARIZED μ^- MESONS. H. Überall, Carnegie Institute of Technology, Pittsburgh, Pennsylvania (Received January 16, 1959).

We investigate the effect of the hyperfine splitting of μ-mesonic atom ground states on the neutron asymmetry from muon capture in hydrogen, in deuterium, and in complex nuclei with spin, which are treated by a one-particle model. It is shown that this can provide more information on the capture interaction than the neutron asymmetry from spinless nuclei. Muon polarizations and gyromagnetic ratios in the hyperfine states are also discussed.

FIXED-ANGLE DISPERSION RELATIONS FOR NUCLEON-NUCLEON SCATTERING. M. Cini, Istituto di Fisica e Scuola di Perfezionamento in Fisica Nucleare dell' Università, Roma, Italia, and Istituto Nazionale di Fisica Nucleare, Sezione di Roma, Roma, Italia, and S. Fubini and A. Stanghellini, CERN, Geneva, Switzerland (Received January 16, 1959).

Starting from the same spectral representation of the nucleon-nucleon scattering amplitude used by Chew in his determination of the pion-nucleon coupling constant, dispersion relations in which the scattering angle θ is kept constant are derived. An extrapolation procedure based on these dispersion relations is proposed with the purpose of obtaining an independent determination of the pion-nucleon coupling constant and the possibility of discriminating between the different available sets of phase shifts.

A first determination of the pion-nucleon coupling constant from these fixed-angle dispersion relations is attempted by means of a simplified extrapolation procedure, closely analogous to the effective-range approximation, and a result in good agreement with the currently accepted value is obtained.

AUTHOR INDEX TO VOLUME 2

References with (E) are to Errata

ABRAHAM, M., M. A. H. McCAUSLAND, and F. N. H. ROBINSON. Dynamic nuclear polarization 449

AIROLDI, G., G. L. BACCHELLA, and E. GERMAGNOLI. Quenched-in lattice vacancies in copper 145

ALBURGER, D. E., A. ELWYN, A. GALLMANN, J. V. KANE, S. OFER, and R. E. PIXLEY. Search for an 80-msec activity in light nuclei 110

ALLARD, J. G. (see Etzel, H. W.) 452

ALLEN, J. S. (see Burman, R. L.) 9

ALLEN, T. K., WILLIAM R. BAKER, ROBERT V. PYLE, and JOHN M. WILCOX. Experimental generation of plasma Alfvén waves 383

ALMQVIST, E. (see Litherland, A. E.) 104

ALVAREZ, LUIS W., PHILIPPE EBERHARD, MYRON L. GOOD, WILLIAM GRAZIANO, HAROLD K. TICHO, and STANLEY G. WOJCICKI. Neutral cascade hyperon event 215

AMADO, R. D. Analysis of deuteron stripping experiments 399

AMBEGAOKAR, VINAY and WALTER KOHN. Connection between true effective mass and optical absorption in insulators 385

ANDERSON, C. E. (see Knox, W. J.) 402

ANDERSON, H. L., T. FUJII, R. H. MILLER, and L. TAU. Electron decay of the positive pion 53; 64(A)

ANDROES, G. M. and W. D. KNIGHT. Nuclear magnetic resonance in superconducting tin 386

APPLE, E. F. (see Halsted, R. E.) 420

ARNS, R. G., R. E. SUND, and M. L. WIEDENBECK. Nuclear resonance fluorescence in Mg^{24} 50

ATKINSON, JOHN H., WILMOT N. HESS, VICTOR PEREZ-MENDEZ, and ROGER W. WALLACE. Neutron cross-section measurements at 4.5 Bev 168

BACCHELLA, G. L. (see Airoldi, G.) 145

BAKER, WILLIAM R. (see Allen, T. K.) 383

BALLAM, J. (see Crittenden, R. R.) 121

BARDON, M., D. BERLEY, and L. M. LEDERMAN. Asymmetry parameter in muon decay 56

_____ (see Berley, David) 357

BARGMANN, V., LOUIS MICHEL, and V. L. TELEGDI. Precession of the polarization of particles moving in a homogeneous electromagnetic field 435

BARKAS, WALTER H., NRIPENDRA N. BISWAS, DONALD A. DeLISE, JOHN N. DYER, HARRY H. HECKMAN, and FRANCIS M. SMITH. Interactions of 1.15-Bev/c K^- mesons in emulsion 466

BARNES, R. G., R. D. ENGARDT, and R. A. HULTSCH. Nuclear magnetic resonance determination of activation volume for diffusion in lithium 202

BARSHAY, SAUL. Remark on strong interactions 315; 437(E)

_____ and SHELDON L. GLASHOW. Possible determination of hyperon parities and coupling strengths 371; 437(E)

BARTON, G. Equivalence transformation of the β-decay operator 224

BATCHELOR, R. (see Litherland, A. E.) 104

BATTERMAN, BORIS W. X-ray measurement of the distribution of electrons in iron and copper 47

BEARD, DAVID B. Microwave emission from high-temperature plasmas 81

BECK, P. A. (see Wei, C. T.) 95

BERGERON, C. J., C. G. GRENIER, and J. M. REYNOLDS. Oscillatory Ettinghausen-Nernst effect 40

BERLEY, DAVID, JULIET LEE, and MARCEL BARDON. Upper limit for the decay mode $\mu \to e+\gamma$ 357

_____ (see Bardon, M.) 56

BERNARDES, N. and H. PRIMAKOFF. Theory of solid He^3 290

BERNSTEIN, JEREMY and TA-YOU WU. X-ray yields in μ-mesonic atomic transitions 404

BIONDI, MANFRED A. and M. P. GARFUNKEL. Measurement of the temperature variation of the energy gap in superconducting aluminum 143

_____ (see Chanin, L. M.) 344
BIRMAN, JOSEPH L. Polarization of fluorescence in CdS and ZnS single crystals 157
BISI, A. and L. ZAPPA. Magnetic perturbation on circular polarization of external bremsstrahlung 348
BISWAS, NRIPENDRA N. (see Barkas, Walter H.) 466
BLATT, F. J. (see Kropschot, R. H.) 91
BLOEMBERGEN, N. Solid state infrared quantum counters 84
BODANSKY, D., S. F. ECCLES, G. W. FARWELL, M. E. RICKEY, and P. C. ROBISON. Time-reversal invariance and the inverse reactions $C^{12}+\alpha \rightleftarrows N^{14}+d$ 101
BÖMMEL, H. E. and K. DRANSFELD. Attenuation of hypersonic waves in quartz 298
BOORSE, H. A. Superconducting electronic specific heats, the "exponential law," and the Bardeen, Cooper, Schrieffer theory 391
BORDONI, P. G., M. NUOVO, and L. VERDINI. Activation energy and relaxation spectrum of dislocations in copper 200
BOWERS, V. A. (see Cochran, E. L.) 43
BOYER, K., E. M. LITTLE, W. E. QUINN, G. A. SAWYER, and T. F. STRATTON. Soft x-rays from a magnetically compressed plasma in Scylla 279
BOYLE, W. S. and K. F. RODGERS. De Haas-van Alphen type oscillations in the infrared transmission of bismuth 338
_____ (see Weinreich, Gabriel) 96
BOYNE, H. S. (see Franken, P. A.) 422
BRAID, T. H. (see Burman, R. L.) 9
BRANSDEN, B. H. and R. G. MOORHOUSE. Σ^--Σ^+ mass excess 431
BRAY, P. J. (see Ring, P. J.) 64(E)
BREIT, G. Effects of atomic electrons on p-p and n-p scattering. II 401
_____ (see Hull, M. H., Jr.) 264
BROCKHOUSE, B. N. Diffusive motions in liquids and neutron scattering 287
_____ Lattice vibrations in silicon and germanium 256
BROLLEY, J. E., JR. (see Rosen, L.) 98
BROMLEY, D. A. (see McKenzie, J. M.) 303
BROUT, R. and H. SUHL. Effects of spin-orbit coupling in rare earth metals, and in solutions of rare earth metals 387
BROWNE, M. E. (see Kaplan, D. E.) 454
BRUNNER, J. H., H. J. LEISI, C. F. PERDRISAT, and P. SCHERRER. Monoenergetic positrons in the decay of Bi^{206} 207
BRUSSEL, M. K. (see Hughes, D. J.) 505
BRYANT, C. A. (see Keesom, P. H.) 260
BURBANK, R. D. and R. D. HEIDENREICH. Extended fine structure of K x-ray absorption edges in Perminvar 85
BURCH, DAVID S. (see Smith, Stephen J.) 165
BURKHARDT, J. L. Inversion of paramagnetic resonance lines in irradiated calcite 149
BURMAN, R. L., W. B. HERRMANNSFELDT, J. S. ALLEN, and T. H. BRAID. Electron-neutrino angular correlation in the beta decay of neon-23 9
BURROWES, H. C., D. O. CALDWELL, D. H. FRISCH, D. A. HILL, D. M. RITSON, and R. A. SCHLUTER. K-meson–nucleon total cross sections from 0.6 to 2.0 Bev 117
_____, D. O. CALDWELL, D. H. FRISCH, D. A. HILL, D. M. RITSON, R. A. SCHLUTER, and M. A. WAHLIG. Pion-proton resonance scattering near 900 Mev 119
BURTON, MILTON (see Dreeskamp, Herbert) 45
BUTTON, KENNETH J., LAURA M. ROTH, W. H. KLEINER, SOLOMON ZWERDLING, and BENJAMIN LAX. Fine structure in the Zeeman effect of excitons in germanium 161

CALDWELL, D. O. (see Burrowes, H. C.) 117
_____ (see Burrowes, H. C.) 119
CAPPS, RICHARD H. Many-channel unitarity condition 475
CARRUTHERS, PETER. Electrical and thermal resistivity of dislocations 336
CHANIN, L. M., A. V. PHELPS, and M. A. BIONDI. Measurement of the attachment of slow electrons in oxygen 344
CHASE, C. E. Velocity of sound in liquid helium near the λ point 197
CHENG, C. H. (see Wei, C. T.) 95
CHOYKE, W. J. (see Patrick, Lyle) 48
CINI, M., E. FERRARI, and R. GATTO. Neutron-proton mass difference by dispersion theory 7
COCHRAN, E. L., V. A. BOWERS, S. N. FONER, and C. K. JEN. Multiple trapping sites for hydrogen atoms in solid argon 43
COCHRAN, WILLIAM. Theory of the lattice vibrations of germanium 495
COHEN, L. D. and W. E. STEPHENS. Gamma-ray activation of carbon 263
COHEN, V. W., J. SCHWARTZ, and R.

NOVICK. Nuclear magnetic moment and spin of europium-152*m* 305
COLLINS, S. A., JR., R. L. KYHL, and M. W. P. STRANDBERG. Spin-lattice relaxation from state of negative susceptibility 88
CORENZWIT, C. (see Hein, R. A.) 500
CRANE, H. R. (see Nelson, D. F.) 492
CRAWFORD, FRANK S., JR., MARCELLO CRESTI, ROGER L. DOUGLASS, MYROL L. GOOD, GEORGE R. KALBFLEISCH, and M. LYNN STEVENSON. Three-body decays of K_2^0 and K_1^0 361
_____, MARCELLO CRESTI, ROGER L. DOUGLASS, MYRON L. GOOD, GEORGE R. KALBFLEISCH, M. LYNN STEVENSON, and HAROLD K. TICHO. Evidence for the $\Delta I = \frac{1}{2}$ rule 266
_____, MARCELLO CRESTI, MYRON L. GOOD, FRANK T. SOLMITZ, and M. LYNN STEVENSON. Further search for parity nonconservation in associated production 11
_____, MARCELLO CRESTI, MYRON L. GOOD, FRANK T. SOLMITZ, M. LYNN STEVENSON, and HAROLD K. TICHO. Λ interactions in hydrogen 174
_____, MARCELLO CRESTI, MYRON L. GOOD, M. LYNN STEVENSON, and HAROLD K. TICHO. Experimental determination of the Λ spin 114
_____, MARCELLO CRESTI, MYRON L. GOOD, M. LYNN STEVENSON, and HAROLD K. TICHO. K^0-K^+ mass excess 112
CRESTI, MARCELLO (see Crawford, Frank S., Jr.) 11
_____ (see Crawford, Frank S., Jr.) 112
_____ (see Crawford, Frank S., Jr.) 114
_____ (see Crawford, Frank S., Jr.) 174
_____ (see Crawford, Frank S., Jr.) 266
_____ (see Crawford, Frank S., Jr.) 361
CREWE, A. V., E. GARWIN, B. LEDLEY, E. LILLETHUN, R. MARCH, and S. MARCOWITZ. Charge independence in the reactions $p+d \rightarrow \pi^0 + He^3$ and $p+d \rightarrow \pi^+ + H^3$ at 450 Mev 269
CRITTENDEN, R. R., J. H. SCANDRETT, W. D. SHEPHARD, W. D. WALKER, and J. BALLAM. π^--p scattering below 1 Bev 121
CZYŻ, W. (see de Swart, J. J.) 51

DAIL, H. W., JR. (see Galt, J. K.) 292
DALITZ, R. H. and S. F. TUAN. Possible resonant state in pion-hyperon scattering 425
DAVIS, H. F., A. ROBERTS, and T. F. ZIPF. Search for the decay $\mu^+ \rightarrow e^+ + \gamma$ 211
DAY, T. B. and G. A. SNOW. Possible explanation of hyperfragment suppression in K^--d reactions 59
_____, G. A. SNOW, and J. SUCHER. Λ^0 and Σ^0 production from the (Σ^-, d) system 468
DeLISE, DONALD A. (see Barkas, Walter H.) 466
DeMARCO, J. J. (see Weiss, R. J.) 148
de SWART, J. J., W. CZYŻ, and J. SAWICKI. Polarized nucleons from the photodisintegration of the deuteron 51
DIRAC, P. A. M. Energy of the gravitational field 368
DORAN, D. G., G. R. FOWLES, and G. A. PETERSON. Shock-wave compression of aluminum 14(E)
DOUGLASS, ROGER L. (see Crawford, Frank S., Jr.) 266
_____ (see Crawford, Frank S., Jr.) 361
DOYLE, W. T., D. J. E. INGRAM, and M. J. A. SMITH. Detection of colloidal centers in lithium hydride by electron resonance 497
DRABBLE, J. R. and R. D. GROVES. Strain-induced changes in the Seebeck coefficient of *n*-type germanium 451
DRANSFELD, K. (see Bömmel, H. E.) 298
DREESKAMP, HERBERT and MILTON BURTON. Measurement of fast luminescence decay times 45
DYER, JOHN N. (see Barkas, Walter H.) 466

EBERHARD, PHILIPPE, ARTHUR H. ROSENFELD, FRANK T. SOLMITZ, ROBERT D. TRIPP, and MASON B. WATSON. K^--hydrogen charge-exchange scattering 312
_____ (see Alvarez, Luis W.) 215
ECCLES, S. F. (see Bodansky, D.) 101
EDMONDS, D. T. and R. G. PETERSEN. Effective exchange constant in yttrium iron garnet 499
ELWYN, A. (see Alburger, D. E.) 110
ENGARDT, R. D. (see Barnes, R. G.) 202
ETZEL, H. W. and J. G. ALLARD. Color center formation in sodium chloride 452
EVERHART, E. (see Ziemba, F. P.) 299

FALGE, R. L., JR. (see Hein, R. A.) 500
FAN, H. Y. (see Fisher, P.) 456

FARWELL, G. W. (see Bodansky, D.) 101
FAUST, W. R. (see Griem, H. R.) 281
FERRARI, E. (see Cini, M.) 7
FERRELL, R. A. (see Tinkham, M.) 331
FISCHER, C. R. (see Hull, M. H., Jr.) 264
FISHER, P. and H. Y. FAN. Optical and magneto-optical absorption effects of Group III impurities in germanium 456
FONDA, LUCIANO and J. E. RUSSELL. Possible determination of $K\Lambda$ relative parity 57
FONER, S. N. (see Cochran, E. L.) 43
FOWLER, T. K. Optical model of inelastic scattering 14(E)
FOWLES, G. R. (see Doran, D. G.) 14(E)
FOX, J. D. (see Hughes, D. J.) 505
FRANKEN, P. A. and H. S. BOYNE. Effect of nuclear electric dipole moments on nuclear spin relaxation in gases 422
FRAZER, WILLIAM R. and JOSE R. FULCO. Effect of a pion-pion scattering resonance on nucleon structure 365
FREDERICKS, W. J. (see Kelly, R. L.) 389
FRISCH, D. H. (see Burrowes, H. C.) 117
_____ (see Burrowes, H. C.) 119
FUJII, T. (see Anderson, H. L.) 53; 64(A)
FULCO, JOSE R. (see Frazer, William R.) 365

GALLMANN, A. (see Alburger, D. E.) 110
GALONSKY, A., H. B. WILLARD, and T. A. WELTON. S-wave detector of deuteron polarization and 14-Mev polarized-neutron source 349
GALT, J. K., F. R. MERRITT, W. A. YAGER, and H. W. DAIL, JR. Cyclotron resonance effects in zinc 292
GARBER, M. (see Kropschot, R. H.) 91
GARFUNKEL, M. P. (see Biondi, Manfred A.) 143
GARWIN, E. (see Crewe, A. V.) 269
GARWIN, R. L., D. P. HUTCHINSON, S. PENMAN, and G. SHAPIRO. Precise determination of the muon magnetic moment 213; 516(E)
GASIOROWICZ, S. G., D. R. YENNIE, and H. SUURA. Magnitude of renormalization constants 513
GATTO, R. (see Cini, M.) 7
GAVENDA, J. D. (see Morse, R. W.) 250
GELLES, I. L. (see Smith, W. V.) 39
GERMAGNOLI, E. (see Airoldi, G.) 145
GERRITSEN, H. J., S. E. HARRISON, H. R. LEWIS, and J. P. WITTKE. Fine structure, hyperfine structure, and relaxation times of Cr^{3+} in TiO_2 (rutile) 153
GETTNER, M., L. HOLLOWAY, D. KRAUS, K. LANDE, E. LEBOY, W. SELOVE, and R. SIEGEL. Measurement of the π^0 mass, and search for a $\pi_0{}^0$ 471
GLASHOW, SHELDON L. (see Barshay, Saul) 371; 437(E)
GLASSGOLD, A. E. and K. R. GREIDER. Generalized diffraction theory for very-high-energy collisions 169
GLICKSMAN, M. and M. C. STEELE. Plasma pinch effects in indium antimonide 461
GOBELI, G. W. (see MacKay, J. W.) 146
GOOD, MYRON L. (see Alvarez, Luis W.) 215
_____ (see Crawford, Frank S., Jr.) 11
_____ (see Crawford, Frank S., Jr.) 112
_____ (see Crawford, Frank S., Jr.) 114
_____ (see Crawford, Frank S., Jr.) 174
_____ (see Crawford, Frank S., Jr.) 266
_____ (see Crawford, Frank S., Jr.) 361
GOVE, H. E. (see Litherland, A. E.) 104
GRAZIANO, WILLIAM (see Alvarez, Luis W.) 215
GREIDER, K. R. (see Glassgold, A. E.) 169
GRENIER, C. G. (see Bergeron, C. J.) 40
GRIEM, H. R., A. C. KOLB, and W. R. FAUST. Bremsstrahlung from dense plasmas 281
GRODZINS, L. and A. W. SUNYAR. Evidence concerning the spin and parity of Eu^{152m} 307
GROVES, R. D. (see Drabble, J. R.) 451
GUPTA, SURAJ N. Theory of spin-orbit interaction in nuclear forces 124

HALSTED, R. E., E. F. APPLE, and J. S. PRENER. Two-stage optical excitation in sulfide phosphors 420
HARRICK, N. J. Properties of a semiconductor surface as determined from a modified drift mobility experiment 199
HARRIS, E. G. Unstable plasma oscillations in a magnetic field 34
HARRISON, S. E. (see Gerritsen, H. J.) 153
HASKIN, D. M. (see Teucher, M. W.) 313
HEBBORN, J. E. and E. H. SONDHEIMER. Diamagnetism of conduction electrons in metals 150
HECKMAN, HARRY H. (see Barkas, Walter H.) 466
HEIDENREICH, R. D. (see Burbank, R.D.) 85

HEIN, R. A., R. L. FALGE, JR., B. T. MATTHIAS, and C. CORENZWIT. Superconductivity versus ferromagnetism in lanthanum-gadolinium alloys 500
HERMAN, FRANK (see Lampert, Murray A.) 394
HERRMANNSFELDT, W. B. (see Burman, R. L.) 9
HESS, WILMOT N. (see Atkinson, John H.) 168
HILL, D. A. (see Burrowes, H. C.) 117
_____ (see Burrowes, H. C.) 119
HOLLOWAY, L. (see Gettner, M.) 471
HOLSTEIN, T. (see Kjeldaas, T., Jr.) 340
HONES, EDWARD W., JR. Antiparticles in space 509
HUGHES, D. J., M. K. BRUSSEL, J. D. FOX, and R. L. ZIMMERMAN. Strengths of high-energy capture gamma rays in W^{184} 505
_____ (see Palevsky, H.) 258
HULDT, LENNART. Optical method for determining carrier lifetimes in semiconductors 3
HULL, D. and H. M. ROSENBERG. Microscopy at liquid helium temperatures: phase transition in sodium 205
HULL, M. H., JR., K. D. PYATT, JR., C. R. FISCHER, and G. BREIT. Nucleon-nucleon spin-orbit interaction potential 264
HULTSCH, R. A. (see Barnes, R. G.) 202
HUTCHINSON, D. P. (see Garwin, R. L.) 213; 516(E)
HWANG, C. F., T. R. OPHEL, E. H. THORNDIKE, RICHARD WILSON, and N. F. RAMSEY. p-p triple scattering at 143 Mev 310

INGRAM, D. J. E. (see Doyle, W. T.) 497

JACCARINO, V. Direct observation of the hfs of a paramagnetic ion in an antiferromagnet. Nuclear magnetic resonance of Co^{59} in CoF_2 163
JACKSON, J. C. and H. W. WYLD, JR. Negative K meson-nucleon interaction at low energies 355
JACOBSEN, E. H. Piezoelectric production of microwave phonons 249
JEN, C. K. (see Cochran, E. L.) 43
JONES, LAWRENCE W. (see Perl, Martin L.) 116

KALBFLEISCH, GEORGE R. (see Crawford, Frank S., Jr.) 266
_____ (see Crawford, Frank S., Jr.) 361
KANE, J. V. (see Alburger, D. E.) 110
KAPLAN, D. E. and M. E. BROWNE. Electron free precession in paramagnetic free radicals 454
KARPLUS, ROBERT, LEROY KERTH, and THADDEUS KYCIA. K-meson nucleon interaction 510
KEESOM, P. H. and C. A. BRYANT. Contribution of the nucleus to the specific heat of rhenium 260
_____ (see Seidel, G.) 261
KELLY, R. L. and W. J. FREDERICKS. Polarization of photoconductivity excitation bands in CdS single crystals 389
KERMAN, A. K., H. McMANUS, and R. M. THALER. Polarization of nucleons in inelastic scattering from nuclei 172
KERNER, EDWARD H. Dynamical Feynman's theorem 152
KERTH, LEROY (see Karplus, Robert) 510
KESSLER, J. O. and A. R. MOORE. Magnetic susceptibility of photogenerated current carriers in germanium 247
KINOSHITA, TOICHIRO. Radiative corrections to π-e decay 477
_____ and ALBERTO SIRLIN. Radiative decay of the muon 177
KISS, Z. J. and H. L. WELSH. Pressure-induced infrared absorption of mixtures of rare gases 166
KJELDAAS, T., JR. and T. HOLSTEIN. Oscillatory magneto-acoustic effect in metals 340
KLEIN, A. (see Tzoar, N.) 433
KLEINER, W. H. and L. M. ROTH. Deformation potential in germanium from optical absorption lines for exciton formation 334
_____ (see Button, Kenneth J.) 161
KLEY, W. (see Palevsky, H.) 258
_____ (see Pelah, I.) 94
KLICK, CLIFFORD C. Energy-level positions of silver luminescent centers in sulfides 418
KLONTZ, E. E. (see MacKay, J. W.) 146
KNIGHT, W. D. (see Androes, G. M.) 386
KNOX, R. S. Inversion symmetry of the M center 87
KNOX, W. J., A. R. QUINTON, and C. E. ANDERSON. Emission of alpha particles

from a compound system of high angular momentum 402
KOHN, WALTER. Image of the Fermi surface in the vibration spectrum of a metal 393
_____ (see Ambegaokar, Vinay) 385
KOLB, A. C. (see Griem, H. R.) 281
KRAUS, D. (see Gettner, M.) 471
KROPSCHOT, R. H., M. GARBER, and F. J. BLATT. Influence of cold work on the resistivity of dilute copper alloys 91
KYCIA, THADDEUS (see Karplus, Robert) 510
KYHL, R. L. (see Collins, S. A., Jr.) 88

LAMPERT, MURRAY A., FRANK HERMAN, and M. C. STEELE. Role of single phonon emission in low-field breakdown of semiconductors at low temperatures 394
LANDE, K. (see Gettner, M.) 471
LANDOVITZ, L. F. and B. MARGOLIS. Strong-coupling theory of the $S = -1$ hyperons: the Λ^0 as a bound state 318
LANE, A. M., J. E. LYNN, E. MELKONIAN, and E. R. RAE. Possible effects of nuclear shell closures on neutron strength functions 424
LAX, BENJAMIN. (see Button, Kenneth J.) 161
LeBOY, E. (see Gettner, M.) 471
LEDERMAN, L. M. (see Bardon, M.) 56
LEDLEY, B. (see Crewe, A. V.) 269
LEE, JULIET (see Berley, David) 357
LEE, T. D. and F. MOHLING. Possible determination of the helicity of elementary excitations in liquid He II 284
LEFKOWITZ, I. (see Pelah, I.) 94
LEINBACH, HAROLD and GEORGE C. REID. Ionization of the upper atmosphere by low-energy charged particles from a solar flare 61
LEISI, H. J. (see Brunner, J. H.) 207
LEMPICKI, A. Polarization of fluorescence in ZnS and CdS single crystals 155
LEWIS, H. R. (see Gerritsen, H. J.) 153
LILLETHUN, E. (see Crewe, A. V.) 269
LIN, C. C. Hydrodynamics of liquid helium II 245
LIPKIN, HARRY J. Collective description of electron interactions without subsidiary conditions 159
LITHERLAND, A. E., E. ALMQVIST, R. BATCHELOR, and H. E. GOVE. Neutron spectra from $Be^9(He^3,n)C^{11}$ and $Be^9(He^4,n)C^{12}$ using pulse shape discrimination in an organic scintillator 104
LITTLE, E. M. (see Boyer, K.) 279
LOHRMANN, E. (see Teucher, M. W.) 313
LONG, DONALD (see Nussbaum, Allen) 6
LOOS, H. G. Heating and confinement of a plasma by a magnetic field of external origin and with a short rise time 282
LOVE, W. A., S. MARDER, I. NADELHAFT, R. T. SIEGEL, and A. E. TAYLOR. Helicity of negative muons from pion decay 107
LOWE, I. J. Free induction decays of rotating solids 285
LUNDY, R. A. (see Swanson, R. A.) 430
LÜTHI, B. Longitudinal magnetoresistance of metals in high fields 503
LYNN, J. E. (see Lane, A. M.) 424

MacARTHUR, J. W. Theory of the origin of the very low-frequency radio-emissions from the earth's exosphere 491
MACFARLANE, G. G., T. P. McLEAN, J. E. QUARRINGTON, and V. ROBERTS. Some effects of strain of thin specimens on absorption measurements at low temperature 252
MacGREGOR, MALCOLM H. p-p S-wave phase shifts from a boundary condition potential 106
MacKAY, J. W., E. E. KLONTZ, and G. W. GOBELI. Thermal and radiation annealing of Ge 146
MADANSKY, L. and G. E. OWEN. Production of polarized proton beams 209
MALLMANN, C. A. System of levels in even-even nuclei 507
MARCH, R. (see Crewe, A. V.) 269
MARCOWITZ, S. (see Crewe, A. V.) 269
MARDER, S. (see Love, W. A.) 107
MARGOLIS, B. (see Landovitz, L. F.) 318
MARSHAK, R. E. (see Okubo, S.) 12
MATTHEWS, P. T. and ABDUS SALAM. Evidence for resonance states in the K^--N system 226
MATTHIAS, B. T. (see Hein, R. L.) 5
_____ (see Suhl, H.) 500
McCAUSLAND, M. A. H. (see Abraham, M.) 449
McKENZIE, J. M. and D. A. BROMLEY. Observation of charged-particle reaction products 303

McLEAN, T. P. (see Macfarlane, G. G.) 252
McMANUS, H. (see Kerman, A. K.) 172
MEISSNER, HANS. Range of order of superconducting electrons 458
MELKONIAN, E. (see Lane, A. M.) 424
MERRITT, F. R. (see Galt, J. K.) 292
MICHEL, LOUIS (see Bargmann, V.) 435
MILLER, R. H. (see Anderson, H. L.) 53; 64(A)
MILLER, ROBERT C. and ALBERT SAVAGE. Direct observation of antiparallel domains during polarization reversal in single-crystal barium titanate 294
MOHLING, F. (see Lee, T. D.) 284
MONTGOMERY, DAVID. Development of hydromagnetic shocks from large-amplitude Alfvén waves 36
MOORE, A. R. (see Kessler, J. O.) 247
MOORHOUSE, R. G. (see Bransden, B. H.) 431
MORAVCSIK, MICHAEL J. Evidence from photoproduction for a pseudoscalar K^+ meson 352
_____ Second resonance in pion photoproduction 171
MORRIS, R. J. (see Siegman, A. E.) 302
MORSE, R. W. and J. D. GAVENDA. Magnetic oscillations of ultrasonic attenuation in a copper crystal at low temperatures 250
MÜLLER, K. ALEXANDER. Electron paramagnetic resonance of manganese IV in $SrTiO_3$ 341
MYERS, JOHN (see Nussbaum, Allen) 6

NADELHAFT, I. (see Love, W. A.) 107
NATHANS, R. and A. PAOLETTI. Magnetic form factor of cobalt 254
NAUENBERG, MICHAEL. Resonances in the π-hyperon system 351
NELSON, D. F., A. A. SCHUPP, R. W. PIDD, and H. R. CRANE. Search for an electric dipole moment of the electron 492
NEUGEBAUER, G., W. D. WALES, and R. L. WALKER. π^-/π^+ ratio from deuterium in photoproduction at 500-1000 Mev 429
NOVICK, R. (see Cohen, V. W.) 305
NUSSBAUM, ALLEN, JOHN MYERS, and DONALD LONG. Effect of hydrostatic pressure on the anomalous Hall coefficient reversal in single crystal tellurium 6

OFER, S. (see Alburger, D. E.) 110
O'KEEFE, J. G. (see Ring, P. J.) 64(E)
OKUBO, S., R. E. MARSHAK, and E. C. G. SUDARSHAN. Isotopic spin selection rules and K_2^0 decay 12
ONEDA, S. and J. C. PATI. V-A four-fermion interaction and the intermediate charged vector meson 125
OPHEL, T. R. (see Hwang, C. F.) 310
OWEN, G. E. (see Madansky, L.) 209

PALEVSKY, H., D. J. HUGHES, W. KLEY, and E. TUNKELO. Lattice vibrations in silicon by scattering of cold neutrons 258
PANOFSKY, W. K. H. and A. N. SAXENA. Search for enhancement of bremsstrahlung produced by 575-Mev electrons in a single crystal of silicon 219
PAOLETTI, A. (see Nathans, R.) 254
PATI, J. C. (see Oneda, S.) 125
PATRICK, LYLE and W. J. CHOYKE. Electron emission from breakdown regions in SiC p-n junctions 48
PELAH, I., I. LEFKOWITZ, W. KLEY, and E. TUNKELO. Observations of hydrogen vibration frequencies in phosphates by means of inelastic scattering of cold neutrons 94
PENMAN, S. (see Garwin, R. L.) 213; 516(E)
PERDRISAT, C. F. (see Brunner, J. H.) 207
PEREZ-MENDEZ, VICTOR (see Atkinson, John H.) 168
PERL, MARTIN L. and LAWRENCE W. JONES. Photography of cosmic rays in a luminescent chamber 116
PETERMANN, A. and Y. YAMAGUCHI. Corrections to the $3D$-$2P$ transitions in μ-mesonic phosphorus and the mass of the muon 359
PETERSEN, R. G. (see Edmonds, D. T.) 499
PETERSON, G. A. (see Doran, D. G.) 14(E)
PHELPS, A. V. (see Chanin, L. M.) 344
PIDD, R. W. (see Nelson, D. F.) 492
PIXLEY, R. E. (see Alburger, D. E.) 110
PONDROM, LEE G. Proposal for a method of polarizing 1-Mev deuterons 346
PRENER, J. S. (see Halsted, R. E.) 420
PRIMAKOFF, H. (see Bernardes, N.) 290
PYATT, K. D., JR. (see Hull, M. H., Jr.) 264
PYLE, ROBERT V. (see Allen, T. K.) 383

QUARRINGTON, J. E. (see Macfarlane, G. G.) 252
QUINN, W. E. (see Boyer, K.) 279

QUINTON, A. F. (see Knox, W. J.) 402

RAE, E. R. (see Lane, A. M.) 424
RAMSEY, N. F. (see Hwang, C. F.) 310
RAPHAEL, R. (see Tzoar, N.) 433
REAGAN, DARYL. Fast ion heating 82
REEVES, HUBERT. Astronomical tests of the existence of Li^4 423
REID, GEORGE C. (see Leinbach, Harold) 61
REYNOLDS, J. M. (see Bergeron, C. J.) 40
RICKAYZEN, G. Collective excitations and the Meissner effect 90
RICKEY, M. E. (see Bodansky, D.) 101
RING, P. J., J. G. O'KEEFE, and P. J. BRAY. Li^7 and F^{19} nuclear magnetic resonances in neutron-irradiated LiF 64(E)
RITSON, D. M. (see Burrowes, H. C.) 117
_____ (see Burrowes, H. C.) 119
ROBERTS, A. (see Davis, H. F.) 211
ROBERTS, V. (see Macfarlane, G. G.) 252
ROBINSON, F. N. H. (see Abraham, M.) 449
ROBISON, P. C. (see Bodansky, D.) 101
RODGERS, K. F. (see Boyle, W. S.) 338
_____ (see Weinreich, Gabriel) 96
ROSEN, L. and J. E. BROLLEY, JR. Time-reversal invariance and parity conservation in strong interactions 98
ROSENBERG, H. M. (see Hull, D.) 205
ROSENFELD, ARTHUR H., FRANK T. SOLMITZ, and ROBERT D. TRIPP $\bar{K}^0$-K^- mass excess 110
_____ (see Eberhard, Philippe) 312
ROTH, LAURA M. (see Button, Kenneth J.) 161
_____ (see Kleiner, W. H.) 334
ROY, P. K. Low-energy application of relativistic K-meson–nucleon dispersion relations 364
RUSSELL, J. E. (see Fonda, Luciano) 57

SALAM, ABDUS. γ_5-symmetrization 127
_____ (see Matthews, P. T.) 226
SAVAGE, ALBERT (see Miller, Robert C.) 294
SAWICKI, J. (see de Swart, J. J.) 51
SAWYER, G. A. (see Boyer, K.) 279
SAXENA, A. N. (see Panofsky, W.K.H.) 219
SCANDRETT, J. H. (see Crittenden, R. R.) 121
SCHERRER, P. (see Brunner, J. H.) 207
SCHIRBER, J. E. and C. A. SWENSON. Superconductivity of β mercury 296
SCHEIN, MARCEL (see Teucher, M. W.) 313
SCHLEIN, PETER E. Nonmesonic/mesonic decay ratio of helium hyperfragments 220
SCHLUTER, R. A. (see Burrowes, H. C.) 117
_____ (see Burrowes, H. C.) 119
SCHNEIDER, JÜRGEN. Stimulated emission of radiation by relativistic electrons in a magnetic field
SCHULZ-DuBOIS, E. O. (see Scovil, H. E. D.) 262
SCHUPP, A. A. (see Nelson, D. F.) 492
SCHWARTZ, J. (see Cohen, V. W.) 305
SCOVIL, H. E. D. and E. O. SCHULZ-DuBOIS. Three-level masers as heat engines 262
SEIDEL, G. and P. H. KEESOM. Nuclear specific heat of gallium and zinc 261
SEIWATZ, HENRY. Measurement of dielectric surface charge density 398
SELOVE, W. (see Gettner, M.) 471
SHAO, PAT. Measurement of dielectric surface charge density and total charge density of a condenser 41
SHAPIRO, G. (see Garwin, R. L.) 213; 516(E)
SHEPHARD, W. D. (see Crittenden, R.R.) 121
SHIREN, N. S. and E. B. TUCKER. Spin-phonon interaction in ruby 206
SHULMAN, R. G. Observation of the Co^{59} nuclear magnetic resonance in paramagnetic salts 459
SIEGEL, R. T. (see Gettner, M.) 471
_____ (see Love, W. A.) 107
SIEGMAN, A. E. and R. J. MORRIS. Proposal for a "staircase" maser 302
SIRLIN, ALBERTO (see Kinoshita, Toichiro) 177
SMITH, FRANCIS M. (see Barkas, Walter H.) 466
SMITH, M. J. A. (see Doyle, W. T.) 497
SMITH, STEPHEN J. and DAVID S. BURCH. Photodetachment cross section of the negative hydrogen ion 165
SMITH, W. V., I. L. GELLES, and P. P. SOROKIN. Electron spin resonance of acceptor states in diamond 39
SNOW, G. A. (see Day, T. B.) 59
_____ (see Day, T. B.) 468
SOLMITZ, FRANK T. (see Crawford, Frank S., Jr.) 11
_____ (see Crawford, Frank S., Jr.) 174
_____ (see Eberhard, Philippe) 312
_____ (see Rosenfeld, Arthur H.) 110
SOLOMON, I. Rotary spin echoes 301
SONDHEIMER, E. H. (see Hebborn, J. E.) 150
SOROKIN, P. P. (see Smith, W. V.) 39
STECH, BERTHOLD. γ_5-invariance and

strong interactions 63
STEELE, M. C. (see Glicksman, M.) 461
_____ (see Lampert, Murray A.) 394
STEIN, P. C. Polarization of recoil protons from the photoproduction of π^0 mesons from hydrogen 473
STEPHENS, W. E. (see Cohen, L. D.) 263
STEVENSON, LYNN (see Crawford, Frank S., Jr.) 11
_____ (see Crawford, Frank S., Jr.) 112
_____ (see Crawford, Frank S., Jr.) 114
_____ (see Crawford, Frank S., Jr.) 174
_____ (see Crawford, Frank S., Jr.) 266
_____ (see Crawford, Frank S., Jr.) 361
STRANDBERG, M. W. P. (see Collins, S. A.) 88
STRATTON, T. F. (see Boyer, K.) 279
STUART, G. W. Satellite-measured radiation 417
SUCHER, J. (see Day, T. B.) 468
SUDARSHAN, E. C. G. (see Okubo, S.) 12
SUHL, H. and B. T. MATTHIAS. Exchange scattering in superconductors 5
_____ (see Brout, R.) 387
SUMI, M. Excitation of oscillations in a plasma layer 37
SUND, R. E. (see Arns, R. G.) 50
SUNYAR, A. W. (see Grodzins, L.) 307
SUURA, H. (see Gasiorowicz, S. G.) 513
SWANSON, R. A., R. A. LUNDY, V. L. TELEGDI, and D. D. YOVANOVITCH. Redetermination of the μ^+ mean life 430
SWENSON, C. A. (see Schirber, J. E.) 296

TAU, L. (see Anderson, H. L.) 53; 64(A)
TAYLOR, A. E. (see Love, W. A.) 107
TELEGDI, V. L. (see Bargmann, V.) 435
_____ (see Swanson, R. A.) 430
TEUCHER, M. W., E. LOHRMANN, D. M. HASKIN, and MARCEL SCHEIN. Nuclear interaction of a proton of about 10^{15} ev producing an electron-photon cascade of 2.4×10^{13} ev 313
THALER, R. M. (see Kerman, A. K.) 172
THORNDIKE, E. H. (see Hwang, C. F.) 310
TICHO, HAROLD K. (see Alvarez, Luis W.) 215
_____ (see Crawford, Frank S., Jr.) 112
_____ (see Crawford, Frank S., Jr.) 114
_____ (see Crawford, Frank S., Jr.) 174
_____ (see Crawford, Frank S., Jr.) 266
TINKHAM, M. and R. A. FERRELL. Determination of the superconducting skin depth from the energy gap and sum rule 331
TRIPP, ROBERT D. (see Eberhard, Philippe) 312
_____ (see Rosenfeld, Arthur H.) 110
TUAN, S. F. (see Dalitz, R. H.) 425
TUCKER, E. B. (see Shiren, N. S.) 206
TUNKELO, E. (see Palevsky, H.) 258
_____ (see Pelah, I.) 94
TZOAR, N., R. RAPHAEL, and A. KLEIN. Meson-theoretical origin of the spin-orbit coupling between two nucleons 433

WAHLIG, M. A. (see Burrowes, H. C.) 119
WALES, W. D. (see Neugebauer, G.) 429
WALKER, R. L. (see Neugebauer, G.) 429
WALKER, W. D. (see Crittenden, R. R.) 121
WALLACE, ROGER W. (see Atkinson, John H.) 168
WATSON, MASON B. (see Eberhard, Philippe) 312
WEI, C. T., C. H. CHENG, and P. A. BECK. Low-temperature specific heats of body-centered cubic Cr-Fe and Cr-Mn alloys 95
WEIBEL, ERICH S. Spontaneously growing transverse waves in a plasma due to an anisotropic velocity distribution 83
WEINBERG, STEVEN. Nature of the vector interaction in μ^--absorption 223
WEINREICH, GABRIEL, W. S. BOYLE, HARRY G. WHITE, and K. F. RODGERS. Valley-orbit splitting of arsenic donor ground state in germanium 96
WEISS, R. J. and J. J. DeMARCO. X-ray measurement of the distribution of electrons in iron and copper 148
WELSH, H. L. (see Kiss, Z. J.) 166
WELTON, T. A. (see Galonsky, A.) 349
WENTZEL, GREGOR. Problem of gauge invariance in the theory of the Meissner effect 33
WHEELER, R. G. Multiplet structure of excitons in CdS 463
WHITE, HARRY G. (see Weinreich, Gabriel) 96
WHITE, ROBERT L. Ferrimagnetic resonance line widths, and g-factors in ferrites 465
WHITFIELD, G. D. Theory of electron-phonon interactions 204
WIEDENBECK, M. L. (see Arns, R. G.) 50

WILCOX, JOHN M. (see Allen, T. K.) 383
WILLARD, H. B. (see Galonsky, A.) 349
WILSON, RICHARD (see Hwang, C. F.) 310
WITTKE, J. P. (see Gerritsen, H. J.) 153
WOJCICKI, STANLEY G. (see Alvarez, Luis W.) 215
WONG, DAVID Y. One-meson contribution to the deuteron quadrupole moment 406
WU, TA-YOU (see Bernstein, Jeremy) 404
WYLD, H. W., Jr. (see Jackson, J. C.) 355

YAGER, W. A. (see Galt, J. K.) 292
YAMAGUCHI, Y. (see Petermann, A.) 359
YENNIE, D. R. (see Gasiorowicz, S. G.) 513
YNTEMA, J. L. and B. ZEIDMAN. Inelastic deuteron scattering from Cu^{63} and Cu^{65} 309
YOVANOVITCH, D. D. (see Swanson, R.A.) 430

ZAPPA, L. (see Bisi, A.) 348
ZEIDMAN, B. (see Yntema, J. L.) 309
ZIEMBA, F. P. and E. EVERHART. Resonance phenomena in large-angle helium ion-helium atom collisions 299
ZIMMERMAN, R. L. (see Hughes, D. J.) 505
ZIPF, T. F. (see Davis, H. F.) 211
ZWERDLING, SOLOMON (see Button, Kenneth J.) 161